# Search for Dark Matter in Association with a Dark Higgs Decaying via two Z Bosons at the ATLAS Detector 



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# Suche nach dunkler Materie in Assoziation mit einem dunklen Higgs-Boson im Zerfallskanal mit zwei Z-Bosonen mit dem ATLAS Detektor 



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#### Abstract

This analysis explores an avenue for testing a theoretical model of dark matter called the Dark Higgs model. This model proposes a majorana fermion $\chi$ as the primary component of the relic dark matter abundance, as well as a $\mathrm{U}(1)$ gauge boson, $\mathrm{Z}^{\prime}$, and a scalar boson, $s$, called the dark Higgs, which couple to the Standard Model. In particular, the objective of the analysis is to optimise signal and control regions for the signal model $s \rightarrow Z Z \rightarrow \ell \bar{\ell} q \bar{q}^{\prime}$, where $\ell$ and $q$ denote kinematically allowed but otherwise arbitrary leptons and quarks respectively. Priority is given to optimising the regions for high dark Higgs masses, as that is where the decay $s \rightarrow \mathrm{ZZ}$ is projected to have a relatively high branching ratio, and other analyses looking at fully hadronic and semileptonic decay modes of $s \rightarrow W W$ already have good sensitivity for low dark Higgs masses. Standard Model background processes resulting from proton-proton collissions at the LHC were simulated with Monte Carlo methods, and recorded with a simulation of the ATLAS detector. Signal processes were also generated with Monte Carlo methods for $Z^{\prime}$ masses ranging from 500 GeV to 3300 GeV , and s masses ranging from 160 GeV to 385 GeV . Different signal and control regions were defined depending on what method was used to reconstruct the hadronically decaying $Z$. The resulting signal and control regions achieve expected exclusion of signal points roughly up to $s$ masses of 400 GeV and down to 170 GeV , for $\mathrm{Z}^{\prime}$ masses between 500 GeV and 2100 GeV at $95 \%$ confidence level. Reconstructing the $s$ with the two jets with the highest transverse momentum achieved the best exclusion for higher $s$ masses, while an algorithm that minimises the invariant mass difference between a pair of jets and the $Z$ boson achieved better exclusion for high $Z^{\prime}$ masses.


To my grandfather, Friðrik Jóhannsson

## Contents

1 Introduction ..... 1
2 Theoretical Background ..... 3
2.1 The Standard Model ..... 3
2.1.1 From Quantum Mechanics to Quantum Field Theory ..... 3
2.1.2 Lagrangian Formulation of QFTs ..... 3
2.1.3 Quantum Electrodynamics ..... 4
2.1.4 The Brout-Englert-Higgs Mechanism and Electroweak Theory ..... 5
2.1.5 The Leptons ..... 7
2.1.6 Quantum Chromodynamics ..... 8
2.2 A Dark Higgs Model ..... 9
3 Experiment ..... 13
3.1 The Large Hadron Collider ..... 13
3.2 The ATLAS Detector ..... 15
3.2.1 Coordinate System ..... 15
3.2.2 Magnet System ..... 16
3.2.3 Inner Detector ..... 17
3.2.4 Calorimeters ..... 18
3.2.5 Muon Spectrometer ..... 19
3.2.6 Trigger System ..... 19
4 Monte Carlo Simulation ..... 21
4.1 Signal Samples ..... 22
4.2 Background Samples ..... 22
5 Object Reconstruction and Definitions ..... 25
5.1 Object Definitions ..... 25
5.1.1 Electrons ..... 25
5.1.2 Muons ..... 25
5.1.3 Small-R Jets $(\mathrm{R}=0.4)$ ..... 26
5.1.4 Track-Assisted-Reclustered Jets ..... 26
5.1.5 Missing Transverse Energy and Significance ..... 27
5.1.6 Overlap Removal ..... 27
6 Analysis ..... 29
6.1 Signal Characteristics ..... 29
6.2 Signal Region Optimization ..... 30
6.2.1 Variables considered ..... 31
6.2.2 Leading Jets Reconstruction ..... 32
6.2.3 Min $\Delta m$ Reconstruction ..... 36
6.2.4 TAR Jet Reconstruction ..... 38
6.3 Control Regions ..... 40
6.4 Statistical Analysis ..... 40
6.5 Expected exclusion ..... 44
6.5.1 Leading jets ..... 44
6.5.2 $\min \Delta m$ reconstruction ..... 46
6.5.3 TAR jet reconstruction ..... 46
7 Conclusion and Outlook ..... 49
Bibliography ..... 57
A $n-1$ Plots ..... 63

## 1 Introduction

The Standard Model of particle physics (SM) represents the current extent of human knowledge about the fundamental forces of nature and the matter contained within it. It is a quantum field theory (QFT) describing the electric, the weak, and the strong interactions, the fermions that comprise the matter we can interact with, and the Higgs field, which is essential to allow the bosons that carry the weak interaction to have masses while keeping the theory renormalizable.

The SM, however, is known not to be a complete description of nature. Gravity, for instance, cannot be renormalizably quantized in the same manner as the other three interactions[1]. Also, more relevantly for this thesis, there is strong evidence that there is an enormous amount of matter in the universe which is not described by any of the particles of the SM[2].
The hypothesis that there is a great deal of matter in the universe that we cannot observe was put forth with good supporting evidence as early as 1884 by Lord Kelvin, who applied the virial theorem to the Galaxy and compared with observed velocities of stars[3]-although there had been previous such suggestions with less rigorous analysis behind them[4]. At first it was assumed that this matter was the same sort of matter as any other astronomical bodies were composed of-massive astrophysical compact halo objects (MACHOs) in modern terminology-except that it was meteoric or nebulous matter that was not observable with the telescopes of the time. As the $20^{\text {th }}$ century progressed and turned into the $21^{\text {st }}$ century, the evidence kept piling up for the existence of this dark matter (DM), the amount of it was narrowed down to be around six times greater than the amount of regular matter[5], and it was realised that only a small fraction of it could be baryonic[6][7][8].

With baryonic matter ruled out as the dominant component of DM, many plausible candidates have been suggested. One candidate is sterile neutrinos[9], a proposed new species of neutrino which does not interact with the electroweak force. Another is the superpartners of SM particles that are predicted by supersymmetry[10], which have have been and are being studied by, for instance, the ATLAS collaboration[11][12]. Yet another is axions, a NambuGoldstone boson which would result from an additional spontaneously broken $U(1)$ symmetry in the SM that couples to the QCD sector[13]. It is also feasible that general relativity could be inaccurate, and that an alternative theory like modified Newtonian dynamics[14] (MOND) could better account for many of the observations on which the DM hypothesis is based, although it has been shown that MOND cannot account for all observations that support the existence of DM[15].

This thesis, however, looks at a particular model of weakly interacting massive particles (WIMPs), called the dark Higgs model. Theoretically, for a particle to freeze out and become a cold relic-a matter distribution with a free streaming length much lower than a protogalaxy-it must have a mass above $\approx 1-100 \mathrm{keV}$, and for the relic to match the observed DM density, a particle with such a mass must self-annihilate with a cross section on the
order of $\sigma v \approx 10^{-26} \mathrm{~cm}^{3} / \mathrm{s}$ where $v$ is the relative velocity between the annihilating particles[4]. This cross section is of a similar scale as those arising from weak interactions, suggesting that this annihilation could be due to a coupling with the weak interaction. The dark Higgs model postulates a fermionic dark matter particle $\chi$, a vector boson $Z^{\prime}$, and a scalar boson $s$. In this analysis, we examine decays of $s$ into two SM $Z$ bosons, one of which decays into two leptons and the other into quarks. This is signature has so far not been explored. Other publically available analyses have examined the decay processes $s \rightarrow b b[16]$ and $s \rightarrow W^{ \pm} W^{\mp} \rightarrow q \bar{q}^{\prime} q \bar{q}^{\prime}[17]$ at the ATLAS detector and $s \rightarrow W W \rightarrow \ell \nu \ell^{\prime} v^{\prime}[18]$ at the CMS detector.

In this thesis, we first introduce the theory of the Standard Model of particle physics, and the dark Higgs model on which this analysis is based. Next, the LHC and ATLAS experiment are described in order to establish the parameters of the Monte Carlo sample and data generation and detection. Thirdly, the essentials of Monte Carlo sample generation are described as well as how the MC backgrounds and signals used in this analysis were created. After that, we describe the objects that are identified and reconstructed at ATLAS which we will be using in this analysis. Finally we move on to presenting the analysis design and the results of the analysis and drawing conclusions and commenting on the analysis.

## 2 Theoretical Background

### 2.1 The Standard Model

### 2.1.1 From Quantum Mechanics to Quantum Field Theory

Special relativity's major departure from classical mechanics is to treat time and space on the same footing. Instead of spatial distances and time intervals between events being separately invariant between inertial reference frames, the invariant quantity is instead the spacetime interval. That is, given events $\left(t_{1}, \boldsymbol{q}_{1}\right)$ and $\left(t_{2}, \boldsymbol{q}_{2}\right)$ in a single reference frame, the spacetime interval

$$
\begin{equation*}
s^{2}=\left(t_{2}-t_{1}\right)^{2}-\left(\boldsymbol{q}_{2}-\boldsymbol{q}_{1}\right)^{2} \tag{2.1}
\end{equation*}
$$

is the same across all reference frames for the same events.
In adapting quantum mechanics to special relativity, the essential problem we face is that quantum mechanics treats time and space completely differently[19]: spatial position is taken to be an operator whereas time is simply taken to be a parameter as in classical mechanics.

There are two ways one could put time and space on equal footing in quantum mechanics. One is to also consider time as an operator like position-an approach taken in string theory, e.g. Ref. [1]. Another is to make position a simple parameter like time-the approach taken in QFT, and hence the SM.

Thus, instead of just operators defined over all of space, we deal with a mapping of spacetime points to operators, or, an operator valued field. More rigorously, we would talk about operator valued distributions on the Schwartz space of Minkowski space[20], but for this introduction such a level of mathematical detail would be overly cumbersome.

### 2.1.2 Lagrangian Formulation of QFTs

In order to study the behaviour of quantum fields, we need to find some equation or equations that they satisfy. Such equations can be obtained by introducing a functional of the quantum fields and their derivatives, called the Lagrangian density-though in the following we will abuse language a bit and simply call this the Lagrangian, as the actual Lagrangian will not feature in this introduction. We then define the action as the spacetime integral of the Lagrangian, and apply the principle of least action.
Denote spacetime coordinates by $x=\left(x^{0}, x^{1}, x^{2}, x^{3}\right)$. Given quantum fields $\phi_{1}(x), \ldots, \phi_{n}(x)$ and a Lagrangian $\mathcal{L}$, the action is given by

$$
\begin{equation*}
\mathcal{S}=\int \mathrm{d}^{4} x \mathcal{L}\left[\phi_{1}(x), \ldots, \phi_{n}(x), \partial_{\mu} \phi_{1}(x), \ldots, \partial_{\mu} \phi_{n}(x)\right] . \tag{2.2}
\end{equation*}
$$

By imposing the requirement that the fields $\phi_{i}$ minimize the action, one can obtain the EulerLagrange equations

$$
\begin{equation*}
\frac{\partial \mathcal{L}}{\partial \phi_{i}}-\partial_{\mu} \frac{\partial \mathcal{L}}{\partial\left(\partial_{\mu} \phi_{i}\right)}=0 \tag{2.3}
\end{equation*}
$$

These are also called the equations of motion for the fields $\phi_{i}$.

### 2.1.3 Quantum Electrodynamics

For a description of quantum electrodynamics, we require a field describing a spin- $1 / 2$ fermion, $\psi$, and a field describing a vector boson, $A_{\mu}$. Let $\gamma^{\mu}$, where $\mu \in\{0,1,2,3\}$, be $4 \times 4$ gamma matrices obeying the anti-commutation relation

$$
\begin{equation*}
\left\{\gamma^{\mu}, \gamma^{v}\right\}=2 \eta^{\mu v} \mathbb{1}_{4} \tag{2.4}
\end{equation*}
$$

Then the field $\bar{\psi}$ is defined as $\bar{\psi}=\psi^{\dagger} \gamma^{0}$. Given a covector $v_{\mu}$, a slash is used as shorthand for its contraction with the gamma matrices: $\psi=\gamma^{\mu} v_{\mu}$.
With the four gamma matrices, we can define a fifth gamma matrix, $\gamma^{5}=i \gamma^{0} \gamma^{1} \gamma^{2} \gamma^{3}$. This matrix is the chirality operator, with eigenvalues $\pm 1$. A spinor $\psi$ is left-handed if $\gamma^{5} \psi=-\psi$ and right handed if $\gamma^{5} \psi=\psi$. With the chirality operator, we can define projectors which project out the left- and right handed components of a spinor, $P_{L}=\frac{1-\gamma^{5}}{2}$ and $P_{R}=\frac{1+\gamma^{5}}{2}$ respectively. Given $A_{\mu}$, we define the electromagnetic field strength as $F_{\mu \nu}=\partial_{\mu} A_{v}-\partial_{\nu} A_{\mu}$. Furthermore, we will define the covariant derivative, $\mathcal{D}_{\mu}=\partial_{\mu}-i q_{e} A_{\mu}$, where $q_{e}$ is the fundamental charge of the spinor field, in this case the fundamental unit of electric charge. The Lagrangian of quantum electrodynamics can then be written as

$$
\begin{equation*}
\mathcal{L}_{Q E D}=i \bar{\psi}\left(\mathscr{D}-m_{e}\right) \psi-\frac{1}{4} F^{\mu v} F_{\mu v} \tag{2.5}
\end{equation*}
$$

This Lagrangian has the property that, for any function $f: \mathbb{R}^{4} \rightarrow \mathbb{R}$ whose second order partial derivatives commute, the gauge transformation

$$
\begin{align*}
\psi & \rightarrow \psi^{\prime}(x)=e^{i f(x)} \psi(x)  \tag{2.6}\\
A_{\mu} \rightarrow A_{\mu}^{\prime}(x) & =A_{\mu}(x)+\frac{1}{e} \partial_{\mu} f(x) \tag{2.7}
\end{align*}
$$

leaves the Lagrangian invariant. This property is called local $U(1)$ gauge redundancy-or gauge invariance, or gauge symmetry, depending on the author's preference-as the action is invariant when the spinor field is acted on by a mapping into the $\mathrm{U}(1)$ Lie group (i.e. complex numbers with an absolute value of 1).

Note that there is a mass term $i m_{e} \bar{\psi} \psi$ for the spinor field in this Lagrangian, but none for the vector field $A_{\mu}$, meaning that the photon is massless. In fact, we would not be able to form a locally $\mathrm{U}(1)$ invariant lagrangian if there was such a mass term for the photon. Based on the gauge redundancy of $A_{\mu}$ it is also commonly called a gauge boson.
The QED Lagrangian gives rise to various Feynman rules, including one vertex for the interaction between photons, commonly denoted by $\gamma$, electrons, denoted by $e^{-}$, and positrons, denoted by $e^{+}$.


Figure 2.1: A generic vertex in quantum electrodynamics

### 2.1.4 The Brout-Englert-Higgs Mechanism and Electroweak Theory

Now consider if, instead of just local $\mathrm{U}(1)$ gauge redundancy, we wanted a locally $\mathrm{U}(1) \times \mathrm{SU}(2)$ gauge redundant Lagrangian. This is in fact the starting point to deriving the electroweak, or Weinberg-Salam Lagrangian. To derive such a theory which is renormalizable and accurately models the weak and electromagnetic interactions, we will not only need four vector gauge bosons, but a complex scalar as well[21-24]. The following exposition is largely a simplified summary of chapter 87 in Ref. [19].

We will denote the gauge boson corresponding to $\mathrm{U}(1)$ by $B_{\mu}$ in this case, with corresponding field strength tensor $B_{\mu \nu}=\partial_{\mu} B_{v}-\partial_{\nu} B_{\mu}$. The corresponding charge is the weak hypercharge $Y$.

Let $\sigma^{a}$, where $a \in 1,2,3$, be the Pauli matrices and we will call $a$ a Lie algebra index, as opposed to a spacetime coordinate index. We could of course instead choose any basis of the Lie algebra $\mathfrak{s u}(2)$ (or a basis multiplied by $i$ if you use the definition of $\mathfrak{s u}(2)$ more common in mathematics), but the Pauli matrices are a convenient choice. We adopt the notation that if the same Lie algebra index appears twice or more often in a product, then it is summed over. Corresponding to the three matrices $\sigma^{a}$, we will need three vector gauge bosons, $W_{\mu}^{a}$, and we define the symbol $W_{\mu}=W_{\mu}^{a} \sigma^{a}$. The corresponding field strength tensor is $W_{\mu \nu}=$ $\partial_{\mu} W_{v}-\partial_{\nu} W_{\mu}-i g\left[W_{\mu}, W_{\nu}\right]$. The corresponding charge is the three component weak isospin I.

Now assume a complex scalar doublet $\phi$, called the Higgs field. We define a covariant derivative for the Higgs field as

$$
\begin{equation*}
\mathcal{D}_{\mu} \phi=\partial_{\mu} \phi-i\left(\frac{g_{2}}{2} W_{\mu}+g_{1} Y B_{\mu}\right) \phi \tag{2.8}
\end{equation*}
$$

where $Y=-\frac{1}{2} \mathbb{1}_{2}$.
Then form the Lagrangian

$$
\begin{equation*}
\mathcal{L}=(\mathcal{D} \phi)^{\dagger} \mathcal{D} \phi+V(\phi)-\frac{1}{4} \operatorname{tr}\left[W^{\mu v} W_{\mu v}\right]-\frac{1}{4} B^{\mu v} B_{\mu v}, \tag{2.9}
\end{equation*}
$$

where $V$ is the potential term

$$
\begin{equation*}
V(\phi)=\frac{1}{4} \lambda\left(\phi^{\dagger} \phi-\frac{v^{2}}{2}\right)^{2} \tag{2.10}
\end{equation*}
$$

which we will refer to as the BEH-potential. Here $\lambda$ is a parameter to be determined and $v$ is real. Thus $\phi$ has a non-zero vacuum expectation value (VEV) such that $\langle 0| \phi^{\dagger} \phi|0\rangle=\frac{v^{2}}{2}$.

With a global gauge transformation, $\phi$ can be chosen such that

$$
\begin{equation*}
\langle 0| \phi(x)|0\rangle=\frac{1}{\sqrt{2}}\binom{v}{0} . \tag{2.11}
\end{equation*}
$$

This choice is called unitary gauge.
We will see that this construction leads to field redefinitions such that we obtain three massive bosons and one massless gauge boson. First note that we can write out

$$
\frac{g_{2}}{2} W_{\mu}^{a} \sigma^{a}+g_{1} B_{\mu} Y=\frac{1}{2}\left(\begin{array}{cc}
g_{2} W_{\mu}^{3}-g_{1} B_{\mu} & g_{2}\left(W_{\mu}^{2}-i W_{\mu}^{1}\right)  \tag{2.12}\\
g_{2}\left(W_{\mu}^{2}+i W_{\mu}^{1}\right) & -\left(g_{2} W_{\mu}^{3}+g_{1} B_{\mu}\right)
\end{array}\right) .
$$

Note that this matrix is self-adjoint. We will define new fields for simplicity: Let $\theta_{W}=$ $\tan ^{-1}\left(g_{1} / g_{2}\right)$, which is commonly known as the Weinberg angle, or the weak mixing angle. We further define $c_{W}=\cos \left(\theta_{W}\right)$ and $s_{W}=\sin \left(\theta_{W}\right)$ and the fields

$$
\begin{align*}
W_{\mu}^{ \pm} & =\frac{1}{\sqrt{2}}\left(W_{\mu}^{1} \mp W_{\mu}^{2}\right)  \tag{2.13}\\
Z_{\mu} & =c_{W} W_{\mu}^{3}-s_{W} B_{\mu}  \tag{2.14}\\
A_{\mu} & =s_{W} W_{\mu}^{3}+c_{W} B_{\mu} . \tag{2.15}
\end{align*}
$$

Then mass terms for the gauge fields $Z_{\mu}$ and $W_{\mu}^{ \pm}$appear in the Lagrangian, and note that the bottom right entry in the matrix does not enter into the Lagrangian:

$$
\frac{1}{8} g_{2}^{2} v^{2}\left(\begin{array}{ll}
1 & 0
\end{array}\right)\left(\begin{array}{cc}
\frac{1}{c_{W}} Z_{\mu} & \sqrt{2} W_{\mu}^{+}  \tag{2.16}\\
\sqrt{2} W_{\mu}^{-} & \ldots
\end{array}\right)^{2}\binom{1}{0}=-\frac{1}{2}\left(\frac{g_{2} v}{2 c_{W}}\right)^{2} Z^{\mu} Z_{\mu}-\left(\frac{g_{2} v}{2}\right)^{2} W^{+\mu} W_{\mu}^{-}
$$

This leads to the $W$ and $Z$ masses

$$
\begin{equation*}
M_{Z}=\frac{g_{2} v}{2 c_{W}} \quad M_{W}=\frac{g_{2} v}{2} \tag{2.17}
\end{equation*}
$$

Rewriting the Higgs field as

$$
\begin{equation*}
\phi(x)=\frac{1}{\sqrt{2}}\binom{v+H(x)}{0} \tag{2.18}
\end{equation*}
$$

leads to a scalar Higgs boson with field strength $H(x)$ and mass

$$
\begin{equation*}
M_{H}=\sqrt{\frac{\lambda}{2}} v . \tag{2.19}
\end{equation*}
$$

Note that the BEH-potential has a rotational symmetry in $\phi$-more concretely, any global $\mathrm{SU}(2)$ transformation of $\phi$ leaves the potential and Lagrangian invariant-whereas no such symmetry exists for the physical Higgs boson $H$. This is commonly illustrated by plotting the BEH-potential assuming that the components of $\phi$ are real, as is done in Fig. 2.2. The red curve marking the minima of the potential thus represents the possible ground states where symmetry is broken.
That the rearranged Lagrangian does not have an $\mathrm{SU}(2)$ gauge redundancy in the fields $W_{\mu}^{ \pm}$ and $Z_{\mu}$ tells us that the ground state of its theory, the vacuum, is not $\mathrm{SU}(2)$ invariant, and this is called spontaneous symmetry breaking. The Lagrangian itself, of course, still has SU(2) gauge redundancy in the fields $W_{\mu}^{a}$.


Figure 2.2: A potential of a similar form as the BEH-potential

The $A_{\mu}$ field still has no mass term in the Lagrangian, meaning that the vacuum state has a $\mathrm{U}(1)$ gauge redundancy, which is thus called an unbroken symmetry. In fact, since this is a different gauge redundancy than the $\mathrm{U}(1)$ redundancy included in the $\mathrm{SU}(2) \times \mathrm{U}(1)$ group, it is common to subscript this group with its corresponding conserved charge, the electric charge $Q$. Correspondingly, the $S U(2) \times U(1)$ group is commonly subscripted with the weak isospin and hypercharge, $S U(2)_{I} \times U(1)_{Y}$.

Since the vacuum state is not symmetric under the same $\mathrm{SU}(2) \times \mathrm{U}(1)$ gauge redundancy as the Lagrangian, the weak isospin and hypercharge are not conserved by interactions with the Higgs field. The electric charge, however, is a combination of the isospin and hypercharge which is conserved

$$
\begin{equation*}
Q=I^{3}+\frac{1}{2} Y . \tag{2.20}
\end{equation*}
$$

We will omit working out the full electroweak Lagrangian in unitary gauge, but from it one can deduce that the $W^{+}$boson has an electric charge of +1 , that the $Z$ boson and photon have electric charges of 0 and the $W^{-}$boson has -1 electric charge[19].

### 2.1.5 The Leptons

There are three generations of leptons in the SM with $S U(2)_{I} \times U(1)_{Y}$ invariance. We number the generations as 1 to 3 , and each one consists of a fermion with electric charge -1 , its antifermion with electric charge +1 and a corresponding neutrino, an electrically neutral lefthanded fermion.

$$
\begin{equation*}
1:\binom{v_{e}}{e} \quad 2:\binom{v_{\mu}}{\mu} \quad 3:\binom{v_{\tau}}{\tau} \tag{2.21}
\end{equation*}
$$

The leptons do not have mass terms as such, but instead have Yukawa couplings to the Higgs field, except for the neutrinos. Let $i$ denote the generation index for the leptons, and let $\psi_{L}$
and $\psi_{R}$ stand for for the left-and right-handed components of $\psi$, respectively. These particles have Yukawa interaction terms with the Higgs field which give them an effective mass:

$$
\begin{equation*}
\mathcal{L}_{Y u k}=-y_{i j} \epsilon(v+H)\left(\bar{\psi}_{L}^{i} \psi_{R}^{j}+\bar{\psi}_{R}^{i} \psi_{L}^{j}\right) . \tag{2.22}
\end{equation*}
$$

Although the coefficients $y_{i j}$ could in principle mix generations of leptons, any such choice can be diagonalized in such a way that the leptons' kinetic terms are unaffected. Thus the leptons have masses

$$
\begin{equation*}
m_{i}=\frac{y_{i} v}{\sqrt{2}} . \tag{2.23}
\end{equation*}
$$

An interesting feature of the electroweak interaction is that the bosons $W^{ \pm}$only interact with left-handed fermions. In fact, since right-handed neutrinos do not couple to any measurable fields in the SM, it is an open question whether they exist. Relatively recent experiments have shown that neutrinos have non-zero masses,[25] which opens up such a possibility.

### 2.1.6 Quantum Chromodynamics

Like the leptons, there are three generations of quarks, but each generation contains two flavours of quark transforming as an $\mathrm{SU}(2)$ doublet. Each flavour of quark is an $\mathrm{SU}(3)$ triplet, and the charge corresponding to $\mathrm{SU}(3)$ is called colour charge[26]. The quarks are additionally $\mathrm{U}(1)$ invariant, so their full gauge group is $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)$. Since $\mathrm{SU}(2) \times \mathrm{U}(1)$ is a subgroup of this group, and there is no more particle content in the $S M$, this is the full gauge group of the SM.

The gauge boson corresponding to $\mathrm{SU}(3)$ is the gluon, $G_{\mu}^{a}$. As there are eight generators of $\mathrm{SU}(3)$, e.g. the Gell-Mann matrices divided by $2, \lambda^{a} / 2$, there are eight such vector bosons. We will use the notation $G_{\mu}=G_{\mu}^{a} \lambda^{a}$.
It improves legibility to separate the quark fields by $\operatorname{SU}(2)$ doublet, so $u_{\alpha, L}^{i}$ denotes the $\alpha$ th $\operatorname{SU}(3)$ component of the quark with $+1 / 2$ weak isospin and $+1 / 3$ weak hypercharge in generation $i$, and $d_{\alpha, L}^{i}$ does likewise for quarks with $-1 / 2$ weak isospin, and $u_{\alpha, R^{\prime}}^{i} d_{\alpha, R}^{i}$ are likewise for quarks with 0 isospin but $+4 / 3$ and $-2 / 3$ weak hypercharge respectively. We will call the quarks denoted with a $u$ up-type quarks, and the quarks denoted with a $d$ down-type quarks. Then the Yukawa terms that give the quarks their effective masses are:

$$
\begin{equation*}
\mathcal{L}_{Y u k}=-\frac{1}{\sqrt{2}}(v+H)\left(y_{i j}^{\prime} \bar{d}_{\alpha, L}^{i} d_{\alpha, R}^{j}+y_{i j}^{\prime \prime} \bar{u}_{\alpha, L}^{i} u_{\alpha, R}^{j}+h . c .\right) . \tag{2.24}
\end{equation*}
$$

While both $y^{\prime}$ and $y^{\prime \prime}$ can be diagonalised simultaneously, and thus yield clean and neat mass terms, the charged quark currents that couple to the $W^{ \pm}$bosons will then unavoidably mix quark generations:

$$
\begin{align*}
J^{+\mu} & =\bar{d}_{\alpha, L}^{i}\left(V^{\dagger}\right)_{i j} \gamma^{\mu} u_{\alpha, L}^{j}  \tag{2.25}\\
J^{-\mu} & =\bar{u}_{\alpha, L}^{i} V_{i j} \gamma^{\mu} d_{\alpha, L}^{j} . \tag{2.26}
\end{align*}
$$

Where $V$ is the Cabibbo-Kobayashi-Maskawa (CKM) matrix which describes how strongly the weak interaction mixes quark generations.
The covariant derivative corresponding to the strong interaction is

$$
\begin{equation*}
\mathcal{D}_{\mu}=\partial_{\mu}-i g_{s} G_{\mu}-i g_{I} W_{\mu}-i g_{Y} B_{\mu} . \tag{2.27}
\end{equation*}
$$

Let $q_{i}$ denote a quark in generation $i$. Then the strong interaction is described by the Lagrangian

$$
\begin{equation*}
\mathcal{L}_{Q C D}=-\frac{1}{4} \operatorname{tr}\left[G^{\mu v} G_{\mu \nu}\right]+\bar{q}_{i}(i \not D) q_{i} . \tag{2.28}
\end{equation*}
$$

The charge associated with $\mathrm{SU}(3)$ is given the descriptive name of colour charge, from which we get the name quantum chromodynamics (QCD), because of the phenomenon of quark confinement. As the physical separation between quarks increases, so too does the strong interaction between them. Conversely, at smaller length scales, i.e. at higher energies, the coupling between strongly interacting particles decreases, and so more accurate results about QCD can be obtained at higher energies.

Quark confinement arises from the fact that all finite energy states of quarks and gluons must be invariant under $\operatorname{SU}(3)$, i.e. have a net zero colour charge. This can be visualised as the quarks and gluons having one of the colours red, green, blue, anti-red, anti-green or antiblue (the anti-colours can themselves be interpreted as cyan, magenta and yellow, though this is really stretching the metaphor past usefulness). Then any physical combination of these must have a net-white colour when combined according to regular colour addition. Quark combinations are divided into two categories, mesons and baryons. Mesons are composed of a color-anticolour pair, while baryons are composed of a quark triple with colours adding up to white. Quarks can also add up to white in other ways, such as in pentaquarks, but such states are not known to ocurr in nature and were only confirmed to have been produced in labs in 2019 by the LHCb. Mesons, baryons, and other quark matter is collectively known as hadrons.

This has been a quick medium level overview of the SM, going into some specifics about the SM Lagrangian while mostly skating past some mathematical details like representation theory. For further reading, Burgess and Moore's[27] book is all about the standard model and has a relatively quick but rigorous overview of the SM Lagrangian early on. To wrap up, the particle content of the SM and some of their basic properties are presented in table 2.1

### 2.2 A Dark Higgs Model

In this thesis, we base our analysis on the dark Higgs (DH) model presented in Ref. [29]. Experimental results at the LHC[30] have strongly constrained the parameter space in which the DM particles can obtain their relic abundance from direct annihilation into SM particles[31]. By postulating a different lightest state in the dark sector, these constraints can be significantly relaxed. In this model, in addition to the SM Lagrangian, we add a Majorana fermion $\chi$ which is the stable DM candidate, a vector gauge boson $Z_{\mu}^{\prime}$ with $\mathrm{U}(1)$ gauge redundancy, and a new complex Higgs field through which the $Z^{\prime}$ and $\chi$ gain masses via spontaneous symmetry breaking. While the constraints on the mass of the DM particle are relaxed if the dark Higgs, $s$, is the lightest state in the dark sector, we will also examine dark Higgs masses up to nearly twice the mass of the DM particle.
Like with the familiar SM Higgs, the DH field $S$ has a real VEV $w$, and can be rewritten as

$$
\begin{equation*}
S=\frac{s+w}{\sqrt{2}} \tag{2.29}
\end{equation*}
$$

where $s$ is recognized as the physical DH boson. The renormalizable part of the dark sector

Table 2.1: Names, masses (rounded to three significant digits if known to that precision), electric charges and spin of all elementary particles in the SM.[28] The commonly used symbols for the particles are given in parentheses after the particle names. Empty entries represent repeated values.

|  | particle | mass | electric charge [ $e$ ] | spin |
| :---: | :---: | :---: | :---: | :---: |
| leptons | electron (e) | 511 keV | -1 | 1/2 |
|  | muon ( $\mu$ ) | 106 MeV |  |  |
|  | tauon ( $\tau$ ) | 1780 Gev |  |  |
|  | electron neutrino ( $v_{e}$ ) | $<1.1 \mathrm{eV}$ | 0 |  |
|  | muon neutrino ( $v_{\mu}$ ) | $<0.19 \mathrm{MeV}$ |  |  |
|  | tauon neutrino ( $v_{\tau}$ ) | $<18.2 \mathrm{MeV}$ |  |  |
| quarks | up (u) | 2.16 MeV | $+\frac{2}{3}$ |  |
|  | charm (c) | 1.27 GeV |  |  |
|  | top ( $t$ ) | 173 GeV |  |  |
|  | down (d) | 4.67 MeV | $-\frac{1}{3}$ |  |
|  | strange (s) | 93 MeV |  |  |
|  | bottom/beauty (b) | 4.18 GeV |  |  |
| bosons | Higgs (H) | 125 GeV | 0 | 0 |
|  | photon ( $\gamma$ ) | 0 |  | 1 |
|  | gluon (g) |  |  |  |
|  | $Z^{0}$ | 91.2 GeV |  |  |
|  | $W^{ \pm}$ | 80.4 GeV | $\pm 1$ |  |

Lagrangian is:

$$
\begin{equation*}
\mathcal{L}_{\chi}=-\frac{1}{2} g^{\prime} q_{\chi} Z^{\prime \mu} \bar{\chi} \gamma^{5} \gamma_{\mu} \chi-\frac{y_{\chi}}{2 \sqrt{2}} s \bar{\chi} \chi+\frac{1}{2} g^{\prime 2} q_{S}^{2} Z^{\prime \mu} Z_{\mu}^{\prime}\left(s^{2}+2 s w\right) . \tag{2.30}
\end{equation*}
$$

Here $g^{\prime}$ denotes the dark sector $\mathrm{U}(1)$ gauge coupling and $q_{\chi}=q_{s} / 2$ is required for gauge invariance. The masses of $\chi$ and $Z^{\prime}$ are $m_{\chi}=y_{\chi} w / \sqrt{2}$ and $m_{Z^{\prime}}=2 g^{\prime} q_{\chi} w$, respectively, while the dark Higgs mass, $m_{s}$, is an independent parameter. This leaves $g_{\chi}=g^{\prime} q_{\chi}$ as a coupling parameter, meaning there are four independent parameters in the dark sector. In terms of these parameters, the Lagrangian can be rewritten as:

$$
\begin{equation*}
\mathcal{L}_{\chi}=-\frac{1}{2} g_{\chi} Z^{\prime \mu} \bar{\chi} \gamma^{5} \gamma_{\mu} \chi-g_{\chi} \frac{m_{\chi}}{m_{Z^{\prime}}} s \bar{\chi} \chi+2 g_{\chi} Z^{\prime \mu} Z_{\mu}^{\prime}\left(g_{\chi} s^{2}+m_{Z^{\prime}} s\right) . \tag{2.31}
\end{equation*}
$$

The main interaction between the dark sector and the SM occurs through $Z^{\prime}$ coupling to quarks:

$$
\begin{equation*}
\mathcal{L}_{\chi / S M}=-g_{q} Z^{\prime \mu} \bar{q} \gamma_{\mu} q . \tag{2.32}
\end{equation*}
$$

We also assume a non-zero mixing angle $\theta$ between the DH boson and the SM Higgs boson, so that the DH boson is unstable and can decay into SM particles with negligible lifetime. Since invisible decays of $Z$ bosons are strongly constrained by measurements from LEP,[32] we look for $Z^{\prime}$ masses much greater than $m_{Z}$. The two most contributing Feynman diagrams for the ZZ final state are shown in Figs. 2.3 and 2.4.
In this search we are interested in dark Higgs masses $m_{s}$ greater than $2 m_{Z}$, so that we have a high branching ratio for $s \rightarrow Z Z$. We set four parameters of the theory to

1. $g_{q}=0.25$.


Figure 2.3: s-channel ( DH and DM emitted from the $Z^{\prime}$ )


Figure 2.4: s-channel (DM emitted from $Z^{\prime}$, DH emitted from $D M$ )
2. $g_{\chi}=1$.
3. $m_{\chi}=200 \mathrm{GeV}$.
4. $\theta=0.01$.
and we examine Monte Carlo events produced with values of $m_{s}$ varying from 160 GeV to 385 GeV , and $m_{Z^{\prime}}$ varying from 500 GeV to 3300 GeV .

## 3 Experiment

The European Organization for Nuclear Research (CERN, from the original French name Conseil européen pour la recherche nucléaire) was founded in 1954 by 12 member states with the goal of advancing the frontiers of nuclear physics, and later particle physics. It now has 23 member states, with its headquarters in Geneva. In 2020 it had almost 11400 scientific users of 110 nationalities from institutes in 76 countries[33]. Since its founding, many groundbreaking discoveries in particle physics have been made at CERN, including the $W[34]$ and $Z[35][36]$ bosons, quark-gluon plasma[37], and the Higgs boson[38].
This chapter introduces CERN's current main particle accelerator, the Large Hadron Collider (LHC), as well as the ATLAS experiment, in which the DH search presented in this thesis is embedded.

### 3.1 The Large Hadron Collider

The LHC is the largest and newest particle accelerator at CERN, and capable of acheiving the highest center-of-mass (COM) energies of them all at 14 TeV , although to date it has only run at a COM energy of 13 TeV . The LHC Project was approved in December 1994, initially to be built in two stages, but after securing a substantial commitment from non-Member states in 1996 the CERN Council approved construction in a single stage[40]. It is installed in the tunnel that previously housed the Large Electron-Positron Collider (LEP), with a circumference of 27 km and passing through the Franco-Swiss border.

The LHC's main operating mode is proton-proton ( $p p$ ) collisions, although it was also designed to be able to accelerate and collide heavy ions such as lead ions. The operating mode considered in this thesis is $p p$ collisions. The LEP, being a particle-antiparticle collider, only needed one storage ring, while the LHC is a particle-particle collider which requires two rings with counter-rotating beams. Because of this context, the tunnel in which the LHC is housed is only 3.7 m in diameter, which is too narrow for two separate proton rings. Rather than enlarge the tunnel, the LHC was built with a twin bore design-originally proposed by John Blewett at the Brookhaven National Laboratory-with two sets of superconducting coils and beam channels in a single magnetic and mechanical structure and cryostat. The disadvantage of this design is that the rings are magnetically coupled, which reduces the LHC's flexibility[41].
The LHC is one endpoint of the CERN accelerator complex, illustrated in Fig. 3.1. To begin with, a duoplasmotron strips the electrons from the protons of a hydrogen gas with a very strong electric field. The proton source is then fed into Linac 4-which replaced Linac 2 in 2020-where the protons are accelerated to 3 MeV by a radio frequency quadrupole, then to 50 MeV by drift tube linacs, then to 100 MeV by coupled-cavity drift tube linacs, and finally

The CERN accelerator complex Complexe des accélérateurs du CERN


LHC - Large Hadron Collider // SPS - Super Proton Synchrotron // PS - Proton Synchrotron // AD - Antiproton Decelerator // CLEAR - CERN Linear
Electron Accelerator for Research // AWAKE - Advanced WAKefield Experiment // ISOLDE - Isotope Separator OnLine // REX/HIE-ISOLDE - Radioactive EXperiment/High Intensity and Energy ISOLDE // MEDICIS // LEIR - Low Energy lon Ring // LINAC - LINear ACcelerator //
n_TOF - Neutrons Time Of Flight // HiRadMat - High-Radiation to Materials // Neutrino Platform

Figure 3.1: CERN accelerator complex as of 2022[39]
to 160 MeV by Pi-mode structures. Linac 2, in comparison, stopped at 50 MeV . Linac 4 then injects this beam into the Proton Synchrotron Booster, which injects its beam into the Proton Synchrotron, which in turn injects its beam into the Super Proton Synchrotron (SPS), which finally injects the beam into the LHC, accelerating the beam to 2 GeV ( 1.4 GeV until 2018), 26 Gev , and 450 GeV respectively.

The LHC is composed of eight straight segments and eight curved segments. The straight segments serve as interaction points for the beams-and are hence called Points, labeled 1-8 going clockwise-either for the particle detectors housed by the LHC or for the LHC hardware. Four of the Points house the four main detectors at the LHC: ALICE, ATLAS (A Toroidal LHC ApparatuS), CMS, and LHCb. ATLAS and CMS-at Points 1 and 5 respectivelyare general purpose detectors that target a similar phase space, but with different equipment, and can thus cross-check each other's results. The purpose of ATLAS and CMS is to perform high precision measurements of the SM as well as search for physics beyond the SM[42][43]. At Point 2, ALICE is optimized for heavy-ion collisions in order to study quarkgluon plasma[44]. At Point 8 , LHCb specializes in heavy flavour physics, looking for indirect evidence of new physics related to CP violation and rare decays of bottom and charm hadrons[45]. Aside from the four main experiments, there are four smaller ones at the LHC: LHCf[46], MoEDAL[47], TOTEM[48], and FASER[49] which is expected to start taking data in 2022.

The remaining Points contain equipment necessary for the LHC to function. The beams from the SPS are injected into the LHC at Points 2 and 8 . Points 3 and 7 contain most of the collimating system, which cleans the beams and ensures that no stray particles can damage the equipment. The acceleration of the beams takes place at Point 4, where two radio frequency cavities operate at 400 MHz and add 8 MeV to each beam's energy during injection and 16 MeV during coast. Due to this method of acceleration, the protons are gathered in


Semiconductor tracker

Figure 3.2: A cutaway diagram of the whole ATLAS detector[50]
bunches of around $10^{11}$ protons, with a bunch spacing of 25 ns . Point 6 contains the beam dumping system. When the beams need to be dumped out of the LHC for whatever reason, fast-paced kicker magnets deflect the beams into dump absorbers. The absorbers each consist of a graphite core, contained in a steel cylinder, which is itself encased in 750 tonnes of concrete and iron shielding.

The curved segments, or arcs, contain superconducting NbTi coils that function as dipole magnets which keep the beams on a circular path. The coils are cooled down with liquid helium to a temperature of 1.9 K , and produce a field strength of 8.3 T -for reference, a fridge magnet has a field strength of around 0.01 T .

### 3.2 The ATLAS Detector

The ATLAS detector is a general purpose detector, located at Point 1 in a cavern 100 m below ground.[42] It is designed to accommodate a wide variety of precision measurements of the SM, as well as searches for BSM physics. It has a cylindrical shape, with a length of 44 m and a diameter of 25 m , and with the interaction point in its center. The design of the ATLAS detector is briefly summarized in Fig. 3.2, and in this chapter we will go into more detail about its design.

### 3.2.1 Coordinate System

ATLAS uses a right-handed coordinate system with its origin at the interaction point (IP). The $x$-axis points toward the centre of the LHC ring, the $y$-axis points upward to the surface, and the $z$-axis is thus determined by the right-hand rule, pointing along the beamline. The $x-y$ plane is called the transverse plane, and the azimuthal angle $\phi$ of a point in the transverse plane is the angle around the beam axis in the interval $[-\pi, \pi]$. The polar angle $\theta$ is measured from the beam axis in the interval $[0, \pi]$. Given an object with energy $E$ and momentum
component $p_{z}$ along the $z$-axis (longitudinal momentum), the rapidity $v$ is defined as

$$
\begin{equation*}
v=\frac{1}{2} \ln \left(\frac{E+p_{z}}{E-p_{z}}\right)=\tanh ^{-1}\left(\frac{p_{z}}{E}\right) . \tag{3.1}
\end{equation*}
$$

Differences in rapidity are Lorentz invariant quantities. Taking the high-momentum limit of the rapidity motivates the definition of pseudorapidity

$$
\begin{equation*}
\eta=-\ln \tan \left(\frac{\theta}{2}\right), \tag{3.2}
\end{equation*}
$$

where $\theta$ refers to the polar angle of the object's momentum. Unlike the rapidity, the pseudorapidity is independent of the energy and momentum calibration of the object, and it has a one-to-one correspondence with the polar angle $\theta$ through $\tanh (\eta)=\cos (\theta)$. On the other hand, differences in pseudorapidity are not Lorentz invariant, except for massless particles, or approximately in the high momentum limit.
Using the azimuthal angle and the pseudorapidity, we define a sort of angular displacement between two objects, $\Delta R$

$$
\begin{equation*}
\Delta R=\sqrt{(\Delta \eta)^{2}+(\Delta \phi)^{2}} . \tag{3.3}
\end{equation*}
$$

Because the total momentum and energy in the transverse plane is much better known than longitudinally, we often use transverse momentum and energy rather than the full momentum and energy. The transverse momentum is defined as $p_{T}=\sqrt{p_{x}^{2}+p_{y}^{2}}$ and the transverse energy as $E_{T}=\sqrt{m^{2}+p_{T}^{2}}$. The transverse momentum is a very useful physical quantity at ATLAS, since the total transverse momentum of an event is known to be zero to within very tight tolerances, and thus one can apply momentum conservation to it. The longitudinal momentum, in contrast, is generally not zero and can vary substantially.

### 3.2.2 Magnet System

In order to accurately measure the momenta of charged particles, ATLAS applies a strong magnetic field to its interior, causing the charged particles' trajectories to curve. The charge-to-momentum ratio of the particles can be measured from the curvature of the tracks, and hence the momentum once the charge is known.
The magnetic field inside ATLAS is generated by a system of four large superconducting magnets. The system is 22 m in diameter, 26 m in length and has a stored energy of 1.6 GJ . The magnets operate at 4.5 K .
The components of this system are:

- a solenoid aligned on the beam axis which provides a 2 T axial magnetic field to the inner detector.
- a barrel toroid which provides a 0.5 T toroidal magnetic field to the barrel muon detector.
- two end-cap toroids that provide 1 T toroidal magnetic fields to the end-cap muon detectors.

These components, as well as the tile calorimeter are illustrated in Fig. 3.3.


Figure 3.3: A diagram of the ATLAS magnet system showing the solenoid, barrel toroid and end-cap toroids in bright orange and the tile calorimeter modeled as four layers with different magnetic properties.[51]

### 3.2.3 Inner Detector

A diagram of the inner detector is pictured in Fig. 3.4. The inner detector comprises the first layers of instruments past the beam pipe and before the calorimeters. The layers closest to the pipe are the pixel detectors. These are circuits containing semiconducting silicon detectors as well as readout electronics. The pixels have been hardened to withstand the intense radiation emitted in the $p p$ collisions, and they provide high-precision tracking information about the produced particles. The nominal pixel size is $50 \times 400 \mu \mathrm{~m}^{2}$. Initially, there were three layers of pixel detectors starting at a radius of $\approx 5 \mathrm{~cm}$ from the beam pipe, but in 2014 the insertable B-layer (IBL) was added at a radius of 3.3 cm from the beam pipe[53].

Past the pixel detectors, there are four double layers of silicon microstrip trackers (SCT), one of which is radial and the other with a stereo angle of 40 mrad . Thus, they can measure all the coordinates of particles which pass through them. The pixels and the strips are present in the barrel as well as in the end-caps of the detector. They each cover the region $|\eta|<2.5$.
Beyond the SCTs are the transition radiation trackers (TRTs). This is a gaseous detector composed of layers of 4 mm diameter drift tubes, filled with a xenon-based gas mixture. The drift tubes have an aluminium cathode coated on a carbon fiber reinforced polyimide layer, and a gold plated tungsten wire as an anode, and are embedded in polystyrene foils with varying electric permittivities. The drift tubes are triggered by transition radiation stemming from ultra-relativistic particles traversing the foils. The TRT is less accurate than the pixels and SCTs, but it provides a longer track length and improves electron identification, since transition radiation increases with speed, and electrons with the same momentum as muons and tauons will have a much greater speed. There are 73 layers of 144 cm long axially aligned tubes in the barrel region, and 160 layers of radially aligned tubes in each of the end caps.


Figure 3.4: Diagrams of the ATLAS inner detector. Adapted from Ref. [52]

### 3.2.4 Calorimeters

Calorimeters are used to measure the energies of particles by absorbing them, so that the heat they transfer is equivalent to their kinetic energies. This is especially important for neutral particles, since they leave no tracks from which their energies and momenta could be deduced. This is very important for studying jets, as they generally contain neutral hadrons. The LHC uses two types of calorimeters, one that targets electrons and photons-the EM calorimeter-and one that targets hadrons-the hadronic calorimeter. Both types are sampling calorimeters, composed of alternating layers of active and absorbing material. In factexcept for the tile calorimeter-they all use liquid argon as the active detector medium.

## Electromagnetic Calorimeter

The EM calorimeter is composed of three parts, a barrel calorimeter and two end-cap calorimeters. All of them use lead as an absorber and have an accordion geometry, that is, the lead panels are scrunched up like the bellows of an accordion. This shape naturally provides full $\phi$ symmetry without any cracks and a fast extraction of the signal at the rear or the front of the electrodes.[51]

The barrel part covers the pseudorapidity range $|\eta|<1.475$ while at the end-caps the calorimeter has two coaxial wheels that cover the ranges $1.375<|\eta|<2.5$ and $2.5<|\eta|<3.2$.

## Hadronic Calorimeters

There are three hadronic calorimeters. Firstly, there is the tile calorimeter, which covers the region directly outside the barrel EM calorimeter as well as two extended barrels, covering the range $|\eta|<1.7$. It uses polystyrene tiles as an active material and steel as an absorber.[51]

The hadronic end-cap calorimeters have a flat-plate design, use copper as an absorber, and cover a range of $1.5<|\eta|<3.2$.[51]

## Forward Calorimeter

The forward calorimeters (FCal) are in the same cryostats as the end-cap calorimeters but are composed of three layers (FCal1-3). The first layer targets electromagnetically interacting particles and uses copper as an absorbing material, while FCal2 and FCal3 mainly use tungsten. These calorimeters cover the range $3.1<|\eta|<4.9$.[51]

### 3.2.5 Muon Spectrometer

The outermost layer of the ATLAS detector is the muon spectrometer (MS). Its job is to detect charged particles that exited the barrel and end-cap calorimeters and measure their momentum in the range $|\eta|<2.7$. It also triggers on charged particles in the range $|\eta|<2.4$. Although in principle the muon spectrometer works for any charged particles, muons are the only charged particles that consistently travel through both the inner detectors and calorimeters without being absorbed.
In the barrel region, the MS is composed of 3 concentric layers, while its end-caps consist of 3 parallel wheels. The magnetic system also applies a magnetic field to the MS which deflects the muons and allows their momenta to be measured. The MS contains the monitored drift tube (MDT) chambers and the cathode strip chambers (CSCs), two high-precision gaseous detector chambers. To synchronize the measurements of the MDT and the CSC with the events they belong to, the resistive plate chambers in the barrel and thin gap chambers in the end-caps are used as triggers.

### 3.2.6 Trigger System

With a bunch crossing occuring every 25 ns at an LHC run, even if there were only one $p p$ collision per bunch crossing, ATLAS cannot even hope to record every single event. Fortunately, this is actually fine since most events that happen per bunch crossing are uninteresting to the ATLAS collaboration anyway. Thus, ATLAS has a trigger system that causes it to only record events if a certain trigger is fired.
ATLAS's most basic trigger is the hardware based L1, which triggers if there are any electrons, muons, tauons, photons, jets or high missing transverse energy in the event. The L1 trigger takes around $2.5 \mu \mathrm{~s}$ to accept the event, bringing the effective event rate at ATLAS down to 100 kHz .

After L1 comes the High-Level Trigger (HLT), which runs in software on a dedicated computing farm. The HLT reduces the event rate even further to 1 kHz , and any data that passes HLT is written to permanent storage at CERN.

## 4 Monte Carlo Simulation

In order to optimize a signal and control region for this analysis, we need theoretical predictions about the events expected to be recorded at ATLAS. This is accomplished by using samples produced using Monte Carlo (MC) methods. As the name suggests, being named after the famous Monacan casino, these methods use random-or more commonly, pseudorandom-inputs to produce probabilistic estimations of various quantities. A simple example would be randomly producing points in a square with side lengths of 1 , and finding the proportion of dots lying inside a circle with radius 1 to the number of total points produced in order to estimate $\pi$. Using MC samples instead of measured data to prepare an analysis strategy is important to prevent cherry-picking the data and obtaining false-positive results.

Monte Carlo methods are on the one hand used to simulate data predicted to be measured at ATLAS assuming only the SM-called background samples-and on the other hand, expected data assuming the DH model-called signal samples. A detailed description of general MC sample generation is give in Ref. [54].

Because some types of events of interest are exceptionally rare, MC generation is usually done by producing events of a specific hard (with high momentum transfer) subprocess. For the number of events of each final state generated by the different MC generators to be proportional to what is expected to be measured at ATLAS, the events are assigned weights based on their cross section[54],

$$
\begin{equation*}
\sigma=\sum_{a, b} \int_{0}^{1} \mathrm{~d} x_{a} \mathrm{~d} x_{b} \int \mathrm{~d} \phi_{n} f_{a}^{h_{1}}\left(x_{a}, \mu_{F}\right) f_{b}^{h_{2}}\left(x_{b}, \mu_{F}\right) \frac{1}{2 x_{a} x_{b} s}\left|\mathcal{M}_{a b \rightarrow n}\right|\left(\phi_{n} ; \mu_{F}, \mu_{R}\right) \tag{4.1}
\end{equation*}
$$

Here $a$ and $b$ are labels for the two interacting partons of the protons, $h_{1}$ and $h_{2}$ label the protons, $x_{a}$ and $x_{b}$ stand for the fractional momentum transfer between the partons, $n$ stands for the number of particles in the final state, and $\phi_{n}$ stands for their phase space. The function $f_{q}^{h}(x, \mu)$ is a parton distribution function (PDF) describing the probability of finding a parton $q$ with a momentum fraction $x$ if probing the hadron $h$ at the energy scale $\mu$. PDFs cannot practically be computed from theory due to significant non-perturbative effects and are instead based on fits to experimental data, detailed in the LHAPDF library[55]. The quantity $\mathcal{M}_{a b \rightarrow n}$ is the matrix element between the initial and final states and is computed in various ways depending on the MC generator[54]. The quantity $\mu_{R}$ is the QCD renormalisation length scale, and $\mu_{F}$ is the factorisation length scale, below which such a perturbative approach can be used[56].
Additional weights can be applied to an event, like experimental corrections to scale effeciencies in MC to measured data efficiencies, or to correct pile-up. Pile-up refers to the number of $p p$-collisions per bunch crossing, and many MC samples were generated before the full pile-up profile was measured in Run 2 of the LHC.

While the hard subprocess is generally calculated to leading order, the partons entering it and some of the particles leaving it can radiate gluons, which can in turn radiate other gluons or quark-antiquark pairs, and so on. These parton showers are simulated with a step-wise Markov chain[54]. This radiation provides a correction to the hard subprocess, and it can be described perturbatively down to an energy scale of 1 GeV , below which hadronization occurs and phenomenological models are used to simulate the processes.

Some of the hadrons produced during hadronization are themselves unstable and quickly decay into partons that again hadronize. At each step where charged particles are involved, QED effects are taken into consideration, possibly resulting in radiated leptons and photons. Additionally, other partons in the protons may interact with each other, which can produce additional parton showers. This part of the event is called the underlying event. The hadrons resulting from these parton showers often hit the detector in narrow cones, called jets.

When working with the generated MC samples, the events also have to be scaled according to the integrated luminosity of the data being compared to. This analysis uses data from Run 2 of the LHC, during which ATLAS reached an integrated luminosity of $139 \mathrm{fb}^{-1}$.

### 4.1 Signal Samples

Given a QFT Lagrangian, Feynman rules describing the interaction or self-interaction of its fields can be derived either by hand, or using computer software like FeynRules. The resulting Universal FeynRules Output can then be used in MC generators to simulate processes of an arbitrary QFT.
In this analysis MadGraph5_AMC@NLO[57] was used to generate $S \rightarrow$ ZZ signal samples to leading order, with NNPDF 3.0 NLO PDF being used to calculate the hard scattering matrix. The parton shower, hadronization and underlying event were simulated using Pythia8[58] with the A14 tuned parameters list[59].
In this analysis, the free parameters of the model are the $Z^{\prime}$ and $S$ masses. The MC samples were produced with the $Z^{\prime}$ mass $m_{Z^{\prime}}$ ranging from 500 GeV to $m_{Z^{\prime}}=3300 \mathrm{GeV}$, and the $S$ mass $m_{S}$ varying between 160 GeV to 385 GeV , as illustrated in Fig. 4.1. Due to very low cross sections, signal samples were not produced with $m_{S}$ above 260 GeV for $m_{Z^{\prime}}=3300 \mathrm{GeV}$, nor with $m_{S}=385 \mathrm{GeV}$ and $m_{Z^{\prime}}=2900 \mathrm{GeV}$.

### 4.2 Background Samples

Powheg-Box v2[60] was used to generate top quark pair processes, collectively referred to as $t \bar{t}$, single top processes, and Higgs processes. For these processes, Pythia8 was used to simulate the parton showers, hadronization, and underlying event.

Processes involving a $W$ or $Z$ boson and jets were simulated with Sherpa 2.2.11[61], diboson processes were simulated with Sherpa 2.2.2 and 2.2.1, and triboson events were simulated with Sherpa 2.2.2.


Figure 4.1: The combinations of $Z^{\prime}$ and $S$ masses for which MC samples were produced, pictured graphically.

## 5 Object Reconstruction and Definitions

This chapter outlines how data recorded by the ATLAS detector is processed to reconstruct the electrons, muons and jet candidates-collectively called objects-that were produced in an event.

### 5.1 Object Definitions

### 5.1.1 Electrons

Electrons are identified by combining data from the electromagnetic (EM) calorimeter and data taken by the inner detector (ID). To begin with, so-called topo-clusters[62] are identified in the electromagnetic and hadronic calorimeters. This is done by identifying a calorimeter cell which has recorded an energy above a certain threshold-which is then called a seed celland its neighbours. If its neighbours also meet the energy threshold, they themselves become seed cells and their neighbours are inspected for their energy, and so on until a connected cluster is identified. While both the hadronic and electromagnetic calorimeters are used to detect a topo-cluster, only the energy recorded by the electromagnetic calorimeter is used to reconstruct the electron, which reduces contamination from pile-up clusters.[62] The topoclusters are then matched to ID-tracks, which are re-fitted to account for energy losses from bremsstrahlung.
The electrons are then divided into two categories, baseline and signal. The baseline and signal electrons have a minimum transverse energy, $E_{T}$, of 7 GeV , a pseudorapidity satisfying $|\eta|<2.47$, and track to vertex associations $\left|d_{0} / \sigma_{d_{0}}\right|<5$, where $d_{0}$ is the transverse offset of the electron candidate's track to the primary vertex and $\sigma_{d_{0}}$ is its uncertainty, and $\Delta z_{0}^{B L} \sin (\theta)<0.5 \mathrm{~mm}$, where $\Delta z_{0}^{B L}$ is the longitudinal offset of the electron candidate's track to the primary vertex. Both types of electrons must also pass a loose isolation criterion, FCLoose. The baseline electrons must pass a filter, LooseAndBLayerLLH, which sets loose identification requirements for the electron candidate. A signal electron must pass a stricter identification filter, MediumLLH.[63]
Because the baseline electrons are reconstructed with a higher efficiency than signal electrons, they are used to reject events with more than two electrons

### 5.1.2 Muons

Muons are reconstructed primarily from ID-tracks and muon spectrometer data.[64] Both baseline and signal muons have a minimum $E_{T}$ of 7 GeV and track associations $d_{0} / \sigma_{d_{0}}<3$ and $\Delta z_{0}^{B L} \sin (\theta)<0.5 \mathrm{~mm}$. Baseline muons have a pseudorapidity range of $|\eta|<2.7$ and pass
the Loose[64] identification criterion, while signal muons are required to have a tighter pseudorapidity range of $|\eta|<2.5$, corresponding to that covered by the inner detector and muon spectrometer, and pass the Medium identification criterion and TightTrackOnly_VarRad isolation criterion.

Muons can possibly be cosmic, with $\left|d_{0}\right|>0.2 \mathrm{~mm}$ and $\left|z_{0}\right|>1 \mathrm{~mm}$, or badly reconstructed, with $\sigma(q / p) /|q / p|>0.2$, where $q$ is the charge and $p$ is the size of the momentum of the muon. These muons are rejected in the analysis.

### 5.1.3 Small-R Jets $(R=0.4)$

Small-R jets are reconstructed using the anti- $k_{t}$ algorithm.[65] Given input particle flow objects labeled with $i$ and $j$, let $k_{t i}$ denote the transverse momentum of object $i, \Delta v_{i j}=v_{i}-v_{j}$ the difference in rapidities between the objects, and $\Delta \phi_{i j}=\phi_{i}-\phi_{j}$ be the difference in azimuthal angles between them. Then we define a distance measure between the objects, $d_{i j}$ and between object $i$ and the beam, $d_{i B}$, by

$$
\begin{align*}
d_{i j} & =\min \left(k_{t i}^{-2}, k_{t j}^{-2}\right) \frac{\Delta \phi_{i j}^{2}+\Delta y_{i j}^{2}}{R^{2}}  \tag{5.1}\\
d_{i B} & =k_{t i}^{-2} \tag{5.2}
\end{align*}
$$

Hard particles are roughly taken to be the centers of jets. For each hard particle $i$, the algorithm goes through surrounding soft particles $j$, calculating these measures, and if $d_{i j}$ is smaller than $d_{i B}$ and smaller than $2 R$, it merges $j$ into a jet with center at $i$. This is repeated until no input objects satisfying the criteria are left. Generally each jet will be conical, except when two or more hard particles are within $2 R$ of each other, in which case the algorithm uses heuristics to decide the shapes of the jets based on their transverse momenta[65][66].

In this analysis, we use fully calibrated jets[67] with $R=0.4$, and require the jets to have minimum transverse momenta, $p_{T}$, of 20 GeV and pseudorapidity range of $|\eta|<2.5$. Jets are cleaned with the TightBad working point to minimize calorimetric noise and non-collision backgrounds. The Tight working point is used to suppress pile-up jets from other $p p$ interactions.

## $b$-tagging

Jets may originate from $b$-quarks, an especially common feature of $t \bar{t}$ backgrounds. These jets are known as $b$-jets and are tagged using a deep learning algorithm called DL1r with $77 \%$ efficiency.[68] A characteristic feature of these jets is that since $b$-hadrons have a long lifetime, there is a secondary vertex from its decay. In this analysis $b$-jets are partially vetoed, so that only events with either no $b$-jets are used, or events with exactly $2 b$-jets where the $b$-jets are used to reconstruct a Z-boson.

### 5.1.4 Track-Assisted-Reclustered Jets

This analysis examines the possibility of using track-assisted-reclustered (TAR) jets to reconstruct the dark Higgs[69]. In this algorithm, the anti- $k_{t}$ algorithm is run again on $R=0.2$ anti- $k_{t}$ with a larger radius parameter $R$, in this case $R=1.0$. Subjets with a $p_{T}$ fraction of less

Table 5.1: Criteria for rejecting one reconstructed object in favour of another.

| Reject | Against | Criteria |
| :--- | :--- | :--- |
| Electron 1 | Electron 2 | shared track, $p_{T 1}<p_{T 2}$ |
| Muon | Electron | shared track, is calorimeter muon |
| Electron | Muon | shared track |
| Jet | Electron | $\Delta R<0.2$ |
| Electron | Jet | $\Delta R<\min \left(0.4,0.04+10 \mathrm{GeV} / p_{T e}\right)$ |
| Jet | Muon | number of tracks $<3$ and $(\mathrm{ghost}$-associated or $\Delta R<0.2)$ |
| Muon | Jet | $\Delta R<\min \left(0.4,0.04+10 \mathrm{GeV} / p_{T \mu}\right)$ |

than 0.05 of the whole jet are removed from the reclustered jet. The input jets are required to have a minimum $p_{T}$ of 20 GeV , a pseudorapidity range of $|\eta|<2.5$.

### 5.1.5 Missing Transverse Energy and Significance

In this analysis, the missing transverse momentum, $E_{T}^{m i s s}$, of an event is calculated based on the momenta of the baseline electrons, the baseline muons, and $R=0.4$ jets.[70] Generally, $E_{T}^{\text {miss }}$ calculations may account for a number of other particles, such as photons or tauons. The calculation also uses a soft terms calculated from tracks that are not associated with any reconstructed objects. Photons and tauons are not included. If the objects are aggregated under the label "visible," then the missing transverse momentum $\vec{p}_{T}^{m i s s}$ is calculated as

$$
\begin{equation*}
\vec{p}_{T}^{\text {miss }}=-\sum_{i} \vec{p}_{i}^{v i s i b l e} . \tag{5.3}
\end{equation*}
$$

The missing transverse energy is the magnitude of this momentum,

$$
\begin{equation*}
E_{T}^{m i s s}=\left|\vec{p}_{T}^{m i s s}\right| . \tag{5.4}
\end{equation*}
$$

This analysis also uses $E_{T}^{m i s s}$ significance,[71] denoted as $\mathcal{S}$, which is calculated based on the uncertainties on the reconstructed objects used to calculate $E_{T}^{m i s s}$, the soft term and a pile-up correction. Let $\vec{p}_{T}^{i n v}$ stand for the combined transverse momentum of all invisible particles, and $L\left(\vec{p}_{T}^{m i s s} \mid \vec{p}_{T}^{\text {inv }}\right)$ stand for the likelihood function of $\vec{p}_{T}^{\text {inv }}$ given a $\vec{p}_{T}^{m i s s}$. Then the significance $\mathcal{S}$ is calculated as,

### 5.1.6 Overlap Removal

When objects are reconstructed, it can happen that different objects are reconstructed with the same data. To remove this double-counting, an overlap removal is applied, resolving the overlap between electrons, muons, and jets, described in table 5.1.

## 6 Analysis

### 6.1 Signal Characteristics

This thesis examines a signal model where a dark Higgs decays into two $Z$ bosons, one of which decays into two leptons and the other into jets. Since lepton flavour and electric charge is conserved in the Standard Model, the leptons are required to be particle-anti-particle pairs of the same generation. Since neutrinos cannot be detected by ATLAS, and tauons are not used in this analysis, this amounts to requiring that the final state of the event have either a signal electron and signal positron, or a signal muon and signal anti-muon, and no other detected leptons. The Z-boson that decays leptonically will be referred to as the leptonic $Z$, or $Z_{\text {cand }}^{l l}$, and is reconstructed by adding the four-momenta of the two signal leptons.

For the hadronically decaying $Z$-boson, referred to as the hadronic $Z$, three reconstruction methods were examined. The first method is simply taking the two jets with the highest transverse momentum and using the sum of their four momenta to reconstruct the $Z$. This candidate $Z$-boson is called $Z_{\text {cand }}^{j j}$, and the method will be referred to as the leading jets method. The second is an algorithm that loops through pairs of jets to find the pair with invariant mass closest to the $Z$-boson mass. This $Z$-boson candidate is called $Z_{\text {cand }}^{\text {min }}$, and the method will be referred to as the $\min \Delta m$ method.
A potential advantage of using $Z_{c a n d}^{m i n \Delta m}$ over $Z_{c a n d}^{j j}$ is that the two jets with the highest transverse momenta are not necessarily the decay products of the hadronic $Z$, and so this approach should allow restricting the hadronic $Z$ mass more without losing as much signal. The downside is that the algorithm artificially inflates the peak around the $Z$ mass in the background, so that a cut on the mass of $Z_{\text {cand }}^{\text {min } \Delta m}$ will not filter out as much background as a cut on the mass of $Z_{\text {cand }}^{j j}$.
The third method is to use the highest momentum TAR jet, in which case the jet itself is effectively treated as the Z-boson. This candidate will be called $Z_{\text {cand }}^{T A R}$. This method identifies the hadronic $Z$ well when it is highly boosted, in which case the jets resulting from its decay will likely be too close together to be well reconstructed as $R=0.4$ jets, but may be well reconstructed as $R=0.2$ jets, thus giving a well reconstructed TAR jet. Highly boosted Zbosons are expected to be more likely the higher the $Z^{\prime}$ - or $s$-mass. The downside of this method is that highly boosted events are relatively rare in general, so it might not give good sensitivity by itself.
Certain baseline requirements are implemented on the samples:

- $E_{T}^{\text {miss }}>100 \mathrm{GeV}$.
- Significance of $E_{T}^{m i s s}>5$.


Figure 6.1: Diagram of the different reconstruction methods

- 2 signal and baseline electrons, or 2 signal and baseline muons.
- $E_{T}^{\text {miss }}$ trigger passed, or single muon trigger passed, or single electron trigger passed.
- Number of $R=0.4 b$-jets is either 0 or 2 .


### 6.2 Signal Region Optimization

A signal region is a region in phase space with an abundance of signal events relative to background events. A signal region is defined by imposing any number of cuts on the dataset under analysis, and various regions defined in this way could be taken as a signal region. The object of this thesis is to find an optimal signal region for the signal $s \rightarrow Z Z \rightarrow \ell \bar{\ell} q \bar{q}^{\prime}$, where $q$ stands for a hadronic jet, defined as being the region where the signal model could be excluded, assuming only the SM, for-roughly speaking-the most combinations of masses of $Z^{\prime}$ and $s$, in other words, the most signal points. We refer to a signal region's ability to exclude signal points as its sensitivity.

It is, however, only roughly speaking, as this analysis of the DH model is only one of many such analyses. Whereas the $s \rightarrow W W$ process theoretically has an overall higher branching ratio, the $s \rightarrow Z Z$ process is expected to have a fairly high branching ratio for high dark Higgs masses. Other analyses have looked at the processes $s \rightarrow W W[18][17]$ and $s \rightarrow b b[16]$, with good significance for dark Higgs masses below around 250 GeV . This analysis will thus focus on obtaining high sensitivity for high dark Higgs masses, possibly at the cost of sensitivity for lower dark Higgs masses.

Optimization was carried out by iteratively creating plots and $n-1$ plots of kinematic quantities, thus identifying cuts on the quantities that yield the greatest Asimov significance $Z$ for the signal samples, where

$$
\begin{equation*}
Z=\sqrt{2 \ln \left(\frac{(s+b)\left(b+\sigma_{b}^{2}\right)}{b^{2}+(s+b) \sigma_{b}^{2}}\right)-\frac{b^{2}}{\sigma_{b}^{2}} \ln \left(1+\frac{\sigma_{b}^{2} s}{b\left(b+\sigma_{b}^{2}\right)}\right)} \tag{6.1}
\end{equation*}
$$

and $s$ is the number of signal events, $b$ is the number of background events and $\sigma_{b}$ is the statistical uncertainty on the number of background events. An $n-1$ plot of a kinematic quantity is a plot on which all cuts imposed on the signal region are applied, except one specific cut on the quantity itself.

### 6.2.1 Variables considered

Here we explore some of the variables used to define the signal and control regions.
Table 6.1: Overview of symbols used throughout.

| symbol | definition |
| :---: | :---: |
| $\Delta R$ (Jet 1, Jet 2) | $\Delta R$ between the jets used to reconstruct $Z_{\text {cand }}^{j j}$ or $Z_{\text {cand }}^{\text {min }}$ m . |
| $\Delta R($ Lepton 1, Lepton 2$)$ | $\Delta R$ between the leptons use to reconstruct $Z_{\text {cand }}^{l l}$. |
| $\Delta R(X, Y)$ | $\Delta R$ between particles $X$ and $Y$. |
| $H_{T}$ | Sum of the magnitudes of the transverse momenta of all jets. |
| $p_{T}($ Jet 2$)$ | $p_{T}$ of the jet with the second highest transverse momentum. |
| $m(X)$ | Mass of particle X. |
| $m^{j j}$ | Mass of $Z_{\text {cfnd }}^{j j}$. |
| $m^{l l}$ | Mass of $Z_{\text {cand }}^{\text {l }}$. |
| $s_{\text {cand }}^{j j}$ | Candidate dark Higgs reconstructed using $Z_{\text {cand }}^{j j}$. |
| $s_{\text {cand }}^{\text {minn }}$ | Candidate dark Higgs reconstructed using $Z_{\text {cand }}^{\text {min }}$ mm . |
| $s_{\text {cand }}^{\text {IAR }}$ | Candidate dark Higgs reconstructed using a TAR jet. |

$\Delta R(X, Y)$
As explained in chapter 3, the variable $\Delta R(X, Y)$ gives a measure of the angular distance between the particles $X$ and $Y$, being higher the further apart they are, and lower the closer together they are.

Since decay products of a parent particle are generally close together, using the $\Delta R$ between the leptonic and hadronic Zs was generally found to exclude background efficiently, being most efficient for lower $m_{s}$. Fig. 6.2 illustrates the effect of cuts on $\Delta R\left(Z_{\text {cand }}^{l l} Z_{\text {cand }}^{j j}\right)$.
The analysis also considered the $\Delta R$ between the jets composing the hadronic $Z$, and found this to often be efficient for isolating signals. The two are complementary, as for higher $m_{s}$ a $\Delta R$ restriction between the reconstructed $Z$ s does not isolate the signal very well, but in turn a restriction on the $\Delta R$ between the subjets of the hadronic $Z$ was found to isolate the signals more efficiently than for lower $m_{s}$.

The $\Delta R$ between the leptons decaying from the leptonic $Z$ was also considered, and found to help with isolating the signal in one signal region. However, a good leptonic $Z$ is usually already identified very well by restricting $m^{l l}$, so this variable is generally not very useful for isolating the signal.

## $E_{T}^{m i s s}$ and the $E_{T}^{m i s s}$ Significance

These variables were defined in chapter 5. A defining feature of dark matter is that it cannot be detected with known apparatuses. Thus, if dark matter is created in a particle collision, it should show up as missing energy in the event, and since we already know that the total transverse momentum of an event is 0 to a very small margin of error, missing transverse momentum provides a good estimate of this missing energy.
The $E_{T}^{\text {miss }}$ significance correlates with $E_{T}^{\text {miss }}$, generally being higher the higher $E_{T}^{m i s s}$ is. However, it was found to be mostly more efficient at removing background than restricting $E_{T}^{m i s s}$.


Figure 6.2: Example $n-1$ plots of $\Delta R\left(Z_{\text {cand }}^{l l}, Z_{\text {cand }}^{j j}\right)$. The left plot is for a signal region targeting high $m_{s}$ while the right targets low $m_{s}$. In both plots, an upper bound on $\Delta R\left(Z_{c a n d}^{l l}, Z_{c a n d}^{j j}\right)$ is indicated by the vertical dashed line. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.

The $E_{T}^{m i s s}$ and $E_{T}^{m i s s}$ significance tend to skew higher the higher the $m_{s}$ of the signal

## Masses of reconstructed $Z$ candidates

A reconstructed $Z$ candidate should have a mass close to that of the $Z$ boson, 91.1876 GeV . Restricting the reconstructed $Z$ candidate masses to be close to the $Z$ boson mass is generally very effective at isolationg the signals irrespective of the $m_{s}$ or $m_{Z^{\prime}}$ of the signal point. As the leptons decaying from the leptonic $Z$ are generally measured more accurately than the hadronic jets, it is possible to restrict the leptonic $Z$ mass, $m^{l l}$, to a tighter window than the hadronic $Z$ mass. Representative $n-1$ plots of $m^{j j}$ are presented in Fig. 6.4

## Other Variables

In some regions, cuts on other variables were incidentally found to efficiently remove background. Placing a lower bound on $H_{T}$ often removed some $t \bar{t}$ background, although because large ranges of $H_{T}$ were empty after all other cuts were applied, the placement of the bounds cannot be said to be precisely nailed down. In another instance, placing a lower bound on the transverse momentum of the jet with the second most transverse momentum improved significance for one signal region.

### 6.2.2 Leading Jets Reconstruction

In the case where the hadronic $Z$ is reconstructed using the leading jets, defining three distinct signal regions targeting different ranges of dark Higgs mass was found to give increased sensitivity over one or two. Since two jets are needed for this reconstruction, a requirement that the number of $R=0.4$ jets be greater than 1 is used across the board. In the case that there are $2 b$-jets present in an event, they are used to reconstruct $Z_{c a n d}^{j j}$. This was found to result in greater sensitivity than a full veto.

The high $s$ mass signal region, which will be named sRLh, is defined by the cuts list in Ta-


Figure 6.3: Example $n-1$ plots of $E_{T}^{m i s s}$ significance in the upper row and $E_{T}^{m i s s}$ in the lower row. The left plots are for a signal region targeting high $m_{s}$ while the right plots target low $m_{s}$. In all plots, a lower bound is indicated by the vertical dashed line. For the bottom right plot, no cut was imposed. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.


Figure 6.4: Example $n-1$ plots of $m^{j j}$. Both plots target the same signal region, but the left plot places a lower bound on $m^{j j}$ indicated by the vertical line while the right plot places an upper bound indicated by the line. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.


Figure 6.5: Plot of the mass of the dark Higgs candidate reconstructed using the leading jets method with selection criteria targeting high dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot depicts the Azimov significance.
ble 6.2, and illustrated with the plots in Figs A. 1 and A. 2 in the appendix. A plot of $m\left(s_{\text {cand }}^{j j}\right)$ is displayed in Fig. 6.5, showing the effectiveness of the signal region. The names of variables used in the plots and throughout are summarised in Table 6.1.

Table 6.2: Selection criteria for the leading jets reconstruction targeting high $m_{s}$, SRLH.

| variable | cut |
| :---: | :---: |
| Significance of $E_{T}^{\text {miss }}$ | $>16$ |
| $E_{T}^{\text {miss }}$ | $>200 \mathrm{GeV}$ |
| $m^{j j}$ | $>80 \mathrm{GeV}$ and $<106 \mathrm{GeV}$ |
| $\Delta R\left(\mathrm{Z}_{\text {cand }}^{l l}, Z_{\text {cand }}^{j j}\right)$ | $<2$ |
| $m^{l l}$ | $>82 \mathrm{GeV}$ and $<98 \mathrm{GeV}$ |
| $\Delta R($ Lepton 1, Lepton 2$)$ | $<2$ |
| $p_{T}($ Jet 2$)$ | $>40 \mathrm{GeV}$ |
| $\Delta R($ Jet 1, Jet 2$)$ | $<1.4$ |
| $H_{T}$ | $<700 \mathrm{GeV}$ |

The signal region targeting medium $s$ masses, named sRLm, is defined by the cuts list in Table 6.3. We display a plot of $m\left(s_{\text {cand }}^{j j}\right)$ with SRLM applied in Fig. 6.6, and the $n-1$ plots used to define the signal region are illustrated in Figs. A. 3 and A. 4 in the appendix.
The signal region targeting low $s$ masses, named sRLL, is defined by the cuts list in Table 6.4. We display $m\left(s_{c a n d}^{j j}\right)$ with sRLL applied in Fig. 6.7 and the $n-1$ plots used to define the region in Figs. A. 5 and A. 6 in the appendix.


Figure 6.6: Plot of the mass of the dark Higgs candidate reconstructed using the leading jets method with selection criteria targeting medium dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.


Figure 6.7: Plot of the mass of the dark Higgs candidate reconstructed using the leading jets method with selection criteria targeting low dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.

Table 6.3: Selection criteria for the leading jets reconstruction targeting medium $m_{s}$, SRLM.

| variable | cut |
| :---: | :---: |
| Significance of $E_{T}^{\text {miss }}$ | $>16$ |
| $E_{T}^{\text {miss }}$ | $>240 \mathrm{GeV}$ |
| $m^{j j}$ | $>70 \mathrm{GeV}$ and $<110 \mathrm{GeV}$ |
| $\Delta R\left(Z_{c a n d l}^{l l}, Z_{\text {cand }}^{j j}\right)$ | $<1.6$ |
| $m^{l l}$ | $>80 \mathrm{GeV}$ and $<100 \mathrm{GeV}$ |
| $\Delta R($ Jet 1, Jet 2$)$ | $<1.8$ |

Table 6.4: Selection criteria for the leading jets reconstruction targeting low $m_{s}$, SRLL.

| variable | cut |
| :---: | :---: |
| Significance of $E_{T}^{\text {miss }}$ | $>14$ |
| $m^{j j}$ | $>60 \mathrm{GeV}$ and $<120 \mathrm{GeV}$ |
| $\Delta R\left(Z_{\text {cand }}^{l l}, Z_{\text {cand }}^{j j}\right)$ | $<0.4$ |
| $m^{l l}$ | $>80 \mathrm{GeV}$ and $<98 \mathrm{GeV}$ |
| $H_{T}$ | $<400 \mathrm{GeV}$ |

### 6.2.3 Min $\Delta m$ Reconstruction

In the case where the dark Higgs candidate, $s_{\text {cand }}^{\operatorname{min\Delta m}}$, is reconstructed using $Z_{\text {cand }}^{\text {min }}$, using two signal regions targeting high and low $m_{s}$ respectively yielded the best results. Since two jets are needed for this reconstruction, it is required that the number of $R=0.4$ jets be greater than 1 in both regions. In the case that there are $2 b$-jets in an event, they are used to reconstruct $Z_{\text {cand }}^{\text {min }}$.
Here, jet 1 and jet 2 refer to the jets used to reconstruct $Z_{\text {cand }}^{\min \Delta m}$. The signal region targeting high $s$ masses, named sRmH, is defined by the cuts list in Table 6.5. We display $m\left(s_{\text {cand }}^{\min \Delta m}\right)$ with sRLL applied in Fig. 6.8 and illustrate the $n-1$ plots in Figs. A. 7 and A. 8 in the appendix.

The signal region targeting low $s$ masses, named SRML, is defined by the cuts list in Table 6.6. We present $m\left(s_{c a n d}^{\min \Delta m}\right)$ with sRmL applied in Fig. 6.9 and illustrate the $n-1$ plots in Figs. A. 9 and A. 10 in the appendix.

Table 6.5: Selection criteria for the $\min \Delta \mathrm{m}$ reconstruction targeting high $m_{s}$, SRMH

| variable | cut |
| :---: | :---: |
| $m\left(Z_{\text {cund }}^{\text {mind }}\right)$ | $>81 \mathrm{GeV}$ and 100 GeV |
| $\Delta R\left(Z_{\text {cand }}^{l} Z_{\text {cand }}^{\text {min } \Delta m}\right)$ | $<1.8$ |
| $\Delta R($ Jet1, Jet $)$ | $<1.6$ |
| $E_{T}^{\text {miss }}$ | $>220 \mathrm{GeV}$ |
| Significance $\left(E_{T}^{\text {miss }}\right)$ | $>18$ |
| $m^{l l}$ | $>83 \mathrm{GeV}$ and $<96$ |
| $H_{T}$ | $<600 \mathrm{GeV}$ |



Figure 6.8: Plot of the mass of the dark Higgs candidate reconstructed using the $\min \Delta m$ method with selection criteria targeting high dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots presents the Azimov significance.


Figure 6.9: Plot of the mass of the dark Higgs candidate reconstructed using the $\min \Delta m$ method with selection criteria targeting low dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.

Table 6.6: Selection criteria for the $\min \Delta \mathrm{m}$ reconstruction targeting low $m_{s}$, SRML.

| variable | cut |
| :---: | :---: |
| $m\left(Z_{\text {cand }}^{\text {mindm }}\right)$ | $>81 \mathrm{GeV}$ and 109 GeV |
| $\Delta R\left(Z_{\text {cand }}^{l l} Z_{\text {cand }}^{\text {min } \Delta m}\right)$ | $<1.2$ |
| $\Delta R($ Jet1, Jet2 $)$ | $<2$ |
| $E_{T}^{\text {miss }}$ | $>200 \mathrm{GeV}$ |
| Significance $\left(E_{T}^{\text {miss }}\right)$ | $>16$ |
| $m^{l l}$ | $>77 \mathrm{GeV}$ and $<100$ |
| $H_{T}$ | $<550 \mathrm{GeV}$ |

Table 6.7: Selection criteria for the TAR jet reconstruction targeting high $m_{s}$, SRTH.

| variable | cut |
| :---: | :---: |
| m(ARJet $\left._{1, R=1.0}\right)$ | $>76 \mathrm{GeV}$ and 100 GeV |
| $\Delta R\left(\right.$ Z $_{\text {cand }}{ }^{l}$ TARJet $\left._{1, R=1.0}\right)$ | $<2$ |
| $E_{T}^{\text {miss }}$ | $>280 \mathrm{GeV}$ |
| Significance $\left(E_{T}^{\text {miss }}\right)$ | $>17$ |
| $m^{l l}$ | $>80 \mathrm{GeV}$ and $<100$ |
| D2 TARjet $\left._{1, R=1.0}\right)$ | $<1.25$ |

### 6.2.4 TAR Jet Reconstruction

In the case where the dark Higgs candidate, $s_{\text {cand }}^{T A R}$, is reconstructed using $Z_{\text {cand }}^{T A R}$, using two signal regions targeting high and low $m_{s}$ respectively yielded the best results. Since one TAR jet is needed for this reconstruction, it is required that the number of $R=1.0$ TAR jets be greater than 0 in both regions. Additionally, the TAR signal regions use a full $b$-jet veto.

Here, TARJet $_{1, R=1.0}$ denotes the TAR jet with the highest transverse momentum, and, for purposes of calculation, is equivalent to $Z_{\text {cand }}^{T A R}$. The variable $D 2\left(\right.$ TARjet $\left._{1, R=1.0)}\right)$ is a ratio of energy correlation functions[72] calculated from the energies and angular distances between the particles comprising the TAR jet. It describes how likely it is that the jet was created by a QCD process rather than a $Z$ boson, with lower values meaning it was likely created from a $Z$ boson, and higher values indicating it originates from QCD processes.
The signal region targeting high $s$ masses, named sRTH, is defined by the cuts list in Table 6.7 and illustrated in Figs. A. 11 and A. 12 in the appendix. We present $m\left(s_{\text {cand }}^{T A R}\right)$ with SRTH applied in Fig. 6.10.
The signal region targeting low $s$ masses, named SRTL, is defined by the cuts list in Table 6.8 and illustrated in Figs. A. 13 and A.14. We present $m\left(s_{\text {cand }}^{T A R}\right)$ with sRTL applied in Fig. 6.11.

Table 6.8: Selection criteria for the TAR jet reconstruction targeting low $m_{s}$, SRTL.

| variable | cut |
| :---: | :---: |
| $m\left(\right.$ TARJet $\left._{1, R=1.0}\right)$ | $>70 \mathrm{GeV}$ and 110 GeV |
| $\Delta R\left(Z_{\text {cand }}^{l l}\right.$, TARJet $\left._{1, R=1.0}\right)$ | $<1.2$ |
| $E_{T}^{\text {miss }}$ | $>240 \mathrm{GeV}$ |
| Significance $\left.^{\text {in }} E_{T}^{\text {miss }}\right)$ | $>14$ |
| $m^{l l}$ | $>82 \mathrm{GeV}$ and $<100$ |
| $D 2\left(\right.$ TARjet $\left._{1, R=1.0}\right)$ | $<1.6$ |



Figure 6.10: Plot of the mass of the dark Higgs candidate reconstructed using a TAR jet with selection criteria targeting high dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.


Figure 6.11: Plot of the mass of the dark Higgs candidate reconstructed using a TAR jet with selection criteria targeting low dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.

Table 6.9: Selection criteria for the control regions

| reconstruction (name) | variable | cut |
| :--- | :--- | :--- |
|  | $E_{T}^{\text {miss }}$ | $>200 \mathrm{GeV}$ |
| all methods | Sig. of $E_{T}^{\text {miss }}$ | $>14$ |
|  | $b$-jets | vetoed |
|  | $m^{l l}$ | $>80 \mathrm{GeV}$ and $<100 \mathrm{GeV}$ |
| leading jets (CRL) | Number of $R=0.4$ jets <br> $m^{j j}$ | $>1$ |
| $\min \triangle \mathrm{~m}($ CRM $)$ | Number of $R=0.4$ jets <br>  <br> $m\left(Z_{\text {cand }}^{\text {mind }}\right)$ | $>110 \mathrm{GeV}$ or $<70 \mathrm{GeV}$ |
| TAR jets (CRT) | Number of TAR jets <br> $m\left(T A R J e t_{1, R=1.0}\right)$ | $>110 \mathrm{GeV}$ or $<70 \mathrm{GeV}$ |
|  | $>110 \mathrm{GeV}$ or $<70 \mathrm{GeV}$ |  |

### 6.3 Control Regions

A control region is a region of phase space enriched in the background that dominantly contaminates the signal region. In this case, the dominant background across the board is the diboson background. Specifically, the diboson samples where the hard process is scattering into two $Z$ bosons, one of which decays into non-neutrino leptons, and the other into neutrinos, cf. Fig. 6.13. This background has a leptonically decaying $Z$ in common with the signal and the neutrinos create missing transverse momentum, also characteristic of the signal. On the other hand, the hadronic $Z$ candidate reconstructed from this background is only reconstructed using coincidental hadronic jets, which will usually result in a poorly reconstructed $Z$, e.g. with a mass very different from the real $Z$ mass.

By defining a control region, one can utilize experimentally measured data to potentially improve the sensitivity of an analysis. This is done by calculating scale factors that align the MC samples to the data, which can then be extrapolated into the signal region. Although a control region may nominally-i.e. without taking into account systematic uncertainties, as is done in this thesis-decrease a signal region's sensitivity, once systematics are taken into account, they should generally improve sensitivity.
The optimized control regions for the reconstruction strategies are given in Table 6.9. Plots of the reconstructed dark Higgs candidate are displayed in Fig. 6.12

When using data in the control region, one must compare the data and MC backgrounds to ensure that there has not been a mismodelling of the processes, or that a significant background has not been missed. We choose a representative data-MC ratio plot of $m\left(S_{\text {cand }}^{j j}\right)$ (Fig. 6.14) to check for anomalies. The shapes of the data and MC agree fairly well. Up to $m_{s} \approx 1100 \mathrm{GeV}$, the data generally has slightly fewer events than the MC, but this is an expected discrepancy which the control region corrects. There is nothing to suggest that a significant background has been excluded or that there was mismodelling.

### 6.4 Statistical Analysis

With the signal and control regions in hand, we can move onto testing their potential power to prove or disprove the DH model for the various signal points. There are two main models for this, one is based on significance and the other is the $\mathrm{CL}_{s}$ method.


Figure 6.12: Plots of dark Higgs candidate masses with their respective control region cuts applied. Signal samples are presented in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ with units of GeV . For readability and due to low statistics, the triboson and single top backgrounds are omitted.


Figure 6.13: The diboson background broken down, illustrated with plots of the mass of the reconstructed dark Higgs candidate and respective control region cuts applied. The qualifier "low" means that the invariant mass of the two leptons are less than 4 GeV . Signal samples are presented in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV .


Figure 6.14: Plot of the mass of the leading jets dark Higgs candidate with the control region applied, including data. Signal samples are presented in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The single top and triboson backgrounds are omitted due to low statistics and for better readability.

First is the concept of significance and $p$-values. Given a hypothesis $H$ and a measurement $M$, one can can quantify how well the measurement agrees with the hypothesis by calculating the probability, based on $H$, of making a measurement $M^{\prime}$ which deviates equally from $H$ as $M$ does, or more so. This probability is called the $p$-value. The $p$-value can also be converted into a significance by matching it to standard deviations of a gaussian, so that the significance of a measurement with $p$-value $p$ is

$$
\begin{equation*}
Z=\Phi^{-1}(1-p) \tag{6.2}
\end{equation*}
$$

Where $\Phi$ is the quantile of the standard Gaussian. To apply this approach, one must first have a hypothesis from which one can calculate quantifiable predictions with quantifiable uncertainties, and one must decide at what $p$-value $M$ can be considered a disproof of $H$. The $p$-value considered a standard varies between disciplines, but in HEP a significance of at least 5-i.e. $5 \sigma$, i.e. a $p$-value lower than about $2.9 \times 10^{-7}$-to be considered positive discovery of a signal. On the other hand, to exclude a signal hypothesis, one usually requires a $p$-value of at least 0.05 , or a significance of about 1.64.[73] The $p$-value tells us how likely it is to measure a given combined signal and background strength. In fact we could rename $p$ as $p_{s+b}$, and consider also just the probability of making the measurement given only the background model, $p_{b}$.
Namely, one runs into trouble with the $p$-value if the signal strength of a model is very small. In that case, an analysis which is not really sensitive to the signal might exclude it, going simply by $p$-values. For such cases the $\mathrm{CL}_{s}$ method has been developed[74]. With this method, the $\mathrm{CL}_{s}$ is defined as

$$
\begin{equation*}
\mathrm{CL}_{s}=\frac{p_{s+b}}{1-p_{b}} \tag{6.3}
\end{equation*}
$$

The denominator penalises small values of $p_{s+b}$, where the value of $p_{b}$ will be close to 1 , and thus the $\mathrm{CL}_{s}$ will grow. Like with the $p$-value, we say that a signal can be excluded if $\mathrm{CL}_{s} \leq 0.05$.

### 6.5 Expected exclusion

Histfitter[75] was used to calculate the expected $\mathrm{CL}_{s}$ values of each signal point. Its utilities were also used to estimate the expected limit of the signal regions' exclusionary potential.
Using the defined signal and control regions, a shape fit was executed to the mass of the dark Higgs candidate of each reconstruction method. Data was used to enhance the control regions and obtain better overall sensitivity.

### 6.5.1 Leading jets

For the leading jet reconstruction, the signal regions were shape fit to $m\left(s_{\text {cand }}^{j j}\right)$ in bins of 10 GeV as summarised in Table 6.10, and using the control region crL. As the signal regions are not orthogonal-i.e. they have events in common-the binning of each region must not overlap with the others.
The binnings of the regions SRLH and SRLM were also explored with SRLM binned up to 320 Gev and down to 290 GeV , and the binning of sRLH was changed accordingly. The best results were yielded by the binning given in Table 6.10, resulting in the exclusion plot shown in Fig. 6.15.

Table 6.10: Summary of $m\left(s_{c a n d}^{j j}\right)$ shape fit binning.

| signal region | $m\left(s_{\text {cand }}^{j j}\right)$ range |
| :--- | :--- |
| SRLH | $310-420 \mathrm{GeV}$ |
| SRLM | $200-310 \mathrm{GeV}$ |
| SRLL | $150-200 \mathrm{GeV}$ |



Figure 6.15: Exclusion plot with $m\left(s_{\text {cand }}^{j j}\right)$ binned in SRM up to 310 GeV . The dashed contour illustrates where $\mathrm{CL}_{s}$ values lower than 0.05 are expected and the shaded region shows the $1 \sigma$ error band.


Figure 6.16: Exclusion plot with signal regions binned in $m\left(s_{\text {cand }}^{\min \Delta m}\right)$. The dashed contour illustrates where $\mathrm{CL}_{s}$ values lower than 0.05 are expected and the shaded region shows the $1 \sigma$ error band.

### 6.5.2 $\min \Delta m$ reconstruction

The signal regions were shape fit to $m\left(s_{\text {cand }}^{\min \Delta m}\right)$ in bins of 10 GeV , with SRML binned from $150-300 \mathrm{GeV}$, and SRмн binned from $300-410 \mathrm{GeV}$, using the control region CRм.

While the $\min \Delta m$ reconstruction provides impressive exclusion on the lower end of the $m_{s}$ spectrum, and towards the higher end of the $m_{Z^{\prime}}$ spectrum, it does have overall worse $\mathrm{CL}_{s}$ for high $m_{s}$ than the leading jets method.

### 6.5.3 TAR jet reconstruction

The signal regions were shape fit to $m\left(s_{\text {cand }}^{\text {TAR }}\right)$ in bins of 10 GeV , with sRTL binned from 150-290 GeV , and SRTh binned from $290-410 \mathrm{GeV}$, using the control region CRT.

As expected, the TAR jets have quite low overall sensitivity


Figure 6.17: Exclusion plot with signal regions binned in $m\left(s_{\text {cand }}^{T A R}\right)$. The dashed contour illustrates where $\mathrm{CL}_{s}$ values lower than 0.05 are expected and the shaded region shows the $1 \sigma$ error band.

## 7 Conclusion and Outlook

To summarise: in this analysis we produced MC signal and background samples. Relevant objects and information was reconstructed from these samples to aid in the analysis, such as the $Z$ candidates and the DH candidate mass. Signal regions resulting from the different hadronic $Z$ reconstruction methods were optimized and eventually compared against each other. Control regions were developed corresponding to the signal regions, and when used with data improved overall exclusion values.

Seeing as this analysis is complementary to other analyses being done with the dark Higgs model in order to boost sensitivity in the high $m_{s}$ region, the leading jets signal regions seem to be the most promising for use in future analysis. Whereas the fully hadronic analysis obtained expected exclusion reaching up to around $270 \mathrm{GeV}[17]$ and observed exclusion reaching slightly below 250 GeV , the analysis presented here has expected exclusion up to $m_{s}=400 \mathrm{GeV}$. For the parameters used in this analysis, however, trying to extrapolate beyond 400 GeV is not possible, since at and above $m_{s}=400 \mathrm{GeV}$ the process $s \rightarrow \chi \chi$ becomes kinematically available (recall that $m_{\chi}$ was set to 200 GeV ) and decays of the $s$ into weak bosons are expected to drop off dramatically. Combining this analysis with previous ones should boost the overall exclusion limit from the $\approx 250 \mathrm{GeV}$ obtained by the fully hadronic analysis.
One possibility that was not explored is combining signal regions for different reconstruction strategies in the same shapefit. This may potentially combine the leading jets' good high $m_{s}$ exclusion with the $\min \Delta m$ reconstruction's better high $m_{Z^{\prime}}$ exclusion at lower $m_{s}$.

While it is common to veto $b$-jets entirely to eliminate $t \bar{t}$ background, this analysis found that allowing certain events containing $b$-jets-in this case, events with two $b$-jets where the $b$-jets were then used to reconstruct the hadronic $Z$-improved sensitivity. As $n-1$ plots of $E_{T}^{m i s s}$ significance in the appendix show, cuts on $E_{T}^{m i s s}$ significance are quite effective at removing $t \bar{t}$ background for this signal model in lieu of a full $b$-jet veto.

By itself, the TAR reconstruction produced definitively the worst results. Often, a scenario where a $Z$ decays into a TAR, or otherwise large $R$, jet is called a merged topology. This is distinguished from a resolved topology where it decays into two well separated jets. Signal regions corresponding to these topologies are sometimes combined in order to boost overall sensitivity. The optimized TAR signal regions and resolved signal regions in this analysis, however, were not orthogonal, and so cannot be directly combined. This analysis did not explore the possibility of imposing a requirement on either signal region in order to make the signal regions orthogonal. Although this is likely to drop statistics in the affected signal region, being able to combine it with the other might make up for this loss of signal and more, boosting sensitivity.

## List of Figures

2.1 A generic vertex in quantum electrodynamics ..... 5
2.2 A potential of a similar form as the BEH-potential ..... 7
2.3 s-channel (DH and DM emitted from the $Z^{\prime}$ ) ..... 11
2.4 s-channel (DM emitted from $Z^{\prime}$, DH emitted from DM ) ..... 11
3.1 CERN accelerator complex as of 2022[39] ..... 14
3.2 A cutaway diagram of the whole ATLAS detector[50] ..... 15
3.3 A diagram of the ATLAS magnet system showing the solenoid, barrel toroid and end-cap toroids in bright orange and the tile calorimeter modeled as four layers with different magnetic properties.[51] ..... 17
3.4 Diagrams of the ATLAS inner detector. Adapted from Ref. [52] ..... 18
4.1 The combinations of $Z^{\prime}$ and $S$ masses for which MC samples were produced, pictured graphically. ..... 23
6.1 Diagram of the different reconstruction methods ..... 30
6.2 Example $n-1$ plots of $\Delta R\left(Z_{\text {cand }}^{l l}, Z_{\text {cand }}^{j j}\right)$. The left plot is for a signal region tar- geting high $m_{s}$ while the right targets low $m_{s}$. In both plots, an upper bound on $\Delta R\left(Z_{\text {cand }}^{l l}, Z_{\text {cand }}^{j j}\right)$ is indicated by the vertical dashed line. Due to low statis- tics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plot presents the Azimov significance. ..... 32
6.3 Example $n-1$ plots of $E_{T}^{\text {miss }}$ significance in the upper row and $E_{T}^{\text {miss }}$ in the lower row. The left plots are for a signal region targeting high $m_{s}$ while the right plots target low $m_{s}$. In all plots, a lower bound is indicated by the vertical dashed line. For the bottom right plot, no cut was imposed. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lowerplot presents the Azimov significance.33
6.4 Example $n-1$ plots of $m^{j j}$. Both plots target the same signal region, but the left plot places a lower bound on $m^{j j}$ indicated by the vertical line while the right plot places an upper bound indicated by the line. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plot presents the Azimov significance.
6.5 Plot of the mass of the dark Higgs candidate reconstructed using the leading jets method with selection criteria targeting high dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot depicts the Azimov significance.
6.6 Plot of the mass of the dark Higgs candidate reconstructed using the leading jets method with selection criteria targeting medium dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.
6.7 Plot of the mass of the dark Higgs candidate reconstructed using the leading jets method with selection criteria targeting low dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.
6.8 Plot of the mass of the dark Higgs candidate reconstructed using the $\min \Delta m$ method with selection criteria targeting high dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots presents the Azimov significance.
6.9 Plot of the mass of the dark Higgs candidate reconstructed using the min $\Delta m$ method with selection criteria targeting low dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.
6.10 Plot of the mass of the dark Higgs candidate reconstructed using a TAR jet with selection criteria targeting high dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.
6.11 Plot of the mass of the dark Higgs candidate reconstructed using a TAR jet with selection criteria targeting low dark Higgs masses applied. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plot presents the Azimov significance.
6.12 Plots of dark Higgs candidate masses with their respective control region cuts applied. Signal samples are presented in the format ( $m_{s}, m_{Z^{\prime}}$ ) with units of GeV . For readability and due to low statistics, the triboson and single top backgrounds are omitted.
6.13 The diboson background broken down, illustrated with plots of the mass of the reconstructed dark Higgs candidate and respective control region cuts applied. The qualifier "low" means that the invariant mass of the two leptons are less than 4 GeV . Signal samples are presented in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV .
6.14 Plot of the mass of the leading jets dark Higgs candidate with the control region applied, including data. Signal samples are presented in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The single top and triboson backgrounds are omitted due to low statistics and for better readability.
6.15 Exclusion plot with $m\left(s_{\text {cand }}^{j j}\right)$ binned in SRM up to 310 GeV . The dashed contour illustrates where $\mathrm{CL}_{s}$ values lower than 0.05 are expected and the shaded region shows the $1 \sigma$ error band.
6.16 Exclusion plot with signal regions binned in $m\left(s_{\text {cand }}^{\text {min } \Delta m}\right)$. The dashed contour illustrates where $\mathrm{CL}_{s}$ values lower than 0.05 are expected and the shaded region shows the $1 \sigma$ error band.
6.17 Exclusion plot with signal regions binned in $m\left(s_{\text {cand }}^{T A R}\right)$. The dashed contour illustrates where $\mathrm{CL}_{s}$ values lower than 0.05 are expected and the shaded region shows the $1 \sigma$ error band.
A. $1 n-1$ plots targeting high $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.
A. $2 n-1$ plots for high $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.
A. $3 n-1$ plots targeting medium $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.
A. $4 n-1$ plots for medium $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.
A. $5 n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.68
A. $6 n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.
A. 7 Various $n-1$ plots targeting high $m_{s}$. Due to low statistics and for readabil- ity, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance. ..... 70
A. 8 Various $n-1$ plots targeting high $m_{s}$. Due to low statistics and for readabil- ity, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance. ..... 71
A. 9 Various $n-1$ plots targeting low $m_{s}$. Due to low statistics and for readabil- ity, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance. ..... 72
A. 10 Various $n-1$ plots targeting low $m_{s}$ and a plot of $m\left(s_{c a n d}^{\min \Delta m}\right)$. Due to low statis- tics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance. ..... 73
A. 11 Various $n-1$ plots targeting high $m_{s}$. Due to low statistics and for readabil- ity, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance. ..... 74
A. 12 Various $n-1$ plots targeting high $m_{s}$ and a plot of $m\left(s_{c a n d}^{T A R}\right)$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance. ..... 75
A. $13 n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the tri- boson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance. ..... 76
A. $14 n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the tri- boson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance. ..... 77

## List of Tables

2.1 Names, masses (rounded to three significant digits if known to that precision), electric charges and spin of all elementary particles in the SM.[28] The com- monly used symbols for the particles are given in parentheses after the particle names. Empty entries represent repeated values. ..... 10
5.1 Criteria for rejecting one reconstructed object in favour of another ..... 27
6.1 Overview of symbols used throughout. ..... 31
6.2 Selection criteria for the leading jets reconstruction targeting high $m_{s}$, SRLH ..... 34
6.3 Selection criteria for the leading jets reconstruction targeting medium $m_{s}$, SRLM ..... 36
6.4 Selection criteria for the leading jets reconstruction targeting low $m_{s}$, SRLL ..... 36
6.5 Selection criteria for the $\min \Delta \mathrm{m}$ reconstruction targeting high $m_{s}$, SRMH ..... 36
6.6 Selection criteria for the $\min \Delta \mathrm{m}$ reconstruction targeting low $m_{s}$, SRML ..... 38
6.7 Selection criteria for the TAR jet reconstruction targeting high $m_{s}$, SRTH. ..... 38
6.8 Selection criteria for the TAR jet reconstruction targeting low $m_{s}$, SRTL. ..... 38
6.9 Selection criteria for the control regions ..... 40
6.10 Summary of $m\left(s_{c a n d}^{j j}\right)$ shape fit binning. ..... 45

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## A n-1 Plots








Figure A.1: $n-1$ plots targeting high $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.


Figure A.2: $n-1$ plots for high $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.


Figure A.3: $n-1$ plots targeting medium $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.


Figure A.4: $n-1$ plots for medium $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.


Figure A.5: $n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.


Figure A.6: $n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.







Figure A.7: Various $n-1$ plots targeting high $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.


Figure A.8: Various $n-1$ plots targeting high $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.


Figure A.9: Various $n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.


Figure A.10: Various $n-1$ plots targeting low $m_{s}$ and a plot of $m\left(s_{\text {cand }}^{\min \Delta m}\right)$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.







Figure A.11: Various $n-1$ plots targeting high $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.


Figure A.12: Various $n-1$ plots targeting high $m_{s}$ and a plot of $m\left(s_{c a n d}^{T A R}\right)$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.







Figure A.13: $n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format ( $m_{s}, m_{Z^{\prime}}$ ) in units of GeV . The lower plots present the Azimov significance.


Figure A.14: $n-1$ plots targeting low $m_{s}$. Due to low statistics and for readability, the triboson and single top backgrounds are omitted. The signal samples are given in the format $\left(m_{s}, m_{Z^{\prime}}\right)$ in units of GeV . The lower plots present the Azimov significance.

## Selbständigkeitserklärung

Ich versichere hiermit, die vorliegende Arbeit mit dem Titel
Suche nach dunkler Materie in Assoziation mit einem dunklen Higgs-Boson im Zerfallskanal mit zwei Z-Bosonen mit dem ATLAS Detektor
selbständig verfasst zu haben und keine anderen als die angegebenen Quellen und Hilfsmittel verwendet zu haben.

Valtýr Kári Daníelsson

München, den 06. April 2022

